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14. ABSTRACT This report results from a contract tasking P. N. Lebedev Physical Institute as follows: The contractor will investigate how high power HF radio waves interact with collisional plasmas, such as the earth's ionosphere. Specifically, the contractor will predict and measure the formation of field aligned small scale striations; and energization of electrons along with their relationship to excited optical emissions. AFRL/AFOSR & AFRL/VS workers (Dr Carlson and coworkers) will provide existing airglow and incoherent scatter radar (ISR) data, and background conditions for boundary inputs to a theoretical calculation. The contractor will perform quantitative comparisons between the parameter values calculated, and those observed.					
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STUDIES OF PLASMA INSTABILITY PROCESSES EXCITED BY GROUND BASED HIGH POWER HF ("HEATING") FACILITIES

The work is driven by the intersection of two lines of activity, one experimental, the other theoretical, that converge on an importantly new view of the physics of how high power HF radio waves interact with collisional plasmas, such as the earth's ionosphere. Its goal is to seize this opportunity to amalgamate the complementary US and Russian recent findings, to strengthen the foundation of plasma physics underlying future work in the growing field of ionospheric HF heating.

The main new results can be summarized as following:

1. Two sets of observations of suprathermal electrons, produced by the interaction of powerful radio wave with the ionosphere, as seen using the incoherent scatter radar (ISR) technique are presented. The observational data are compared with the theory of multiple acceleration of electrons in the strongly excited resonance region near the reflection point of the powerful radio wave. The structure of the wide perturbed region filled with energetic (10–20) eV electrons is determined. The size of this region along the Earth magnetic field is shown to be 100 km. The full power going to the accelerated particles is determined: it is 6–8 kW or (4–6)% of the entire HF radiated power. The power carried by the suprathermal electrons flux escaping into magnetosphere is of the order 1 kW.
2. The of electrons near the O-wave reflection point by combined action of upper hybrid and Langmuir turbulence is considered. It is demonstrated that the calculations of electron acceleration in Langmuir resonance layer without electron collisions lead to controversial result. Kinetic theory of multiple acceleration of suprathermal electrons is developed. The energy losses of electrons on the excitation of molecular vibration level are taken into account. The model calculations describing the general peculiarities of the suprathermal electron space energy distribution are performed.
3. Irregularities of electron temperature and plasma density in the F-region of the ionosphere are strongly elongated magnetic field due to dramatic difference in transport coefficients along the magnetic field and transverse to it. Their scales along the magnetic field determined previously neglecting plasma drift effects are rather long reaching ten or even several tens kilometers. Such elongated irregularities should be affected by ionospheric drifts perpendicular to magnetic field, but up to date there was no theoretical estimations for this influence. Two main problems are solved: constant drift effects on irregularities caused by the sources fixed in the ionosphere (for example, narrow radio beam), and gradient drift effects on irregularities created by the source moving with plasma.

Publications

1. A. V. Gurevich H. C. Carlson, G. M. Milikh, K. P. Zybin, F. T. Djuth and K. Groves "Suprathermal electrons generated by the interaction of powerful radio wave with the ionosphere", GRL 2000 (accepted for publication).
2. A.V. Gurevich, H.C. Carlson, G.M. Milikh and K.P. Zybin "Optic emission from modified by powerful radio waves ionosphere" (in preparation for publication).
3. A.V. Gurevich, A.N. Karashtin, K.P. Zybin "Drift effects on the scale of field-aligned irregularities" (in preparation for publication).

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1. Experimental Studies of Energetic Electrons in Ionosphere Using ISR Technique and Comparison of Observational Data with Multiple Acceleration Theory

Abstract

Two sets of observations of suprathermal electrons, produced by the interaction of powerful radio wave with the ionosphere, as seen using the incoherent scatter radar (ISR) technique are presented. The observational data are compared with the theory of multiple acceleration of electrons in the strongly excited resonance region near the reflection point of the powerful radio wave. The structure of the wide perturbed region filled with energetic (10–20) eV electrons is determined. The size of this region along the Earth magnetic field is shown to be 100 km. The full power going to the accelerated particles is determined: it is 6–8 kW or (4–6)% of the entire HF radiated power. The power carried by the suprathermal electrons flux escaping into magnetosphere is of the order 1 kW.

1.1 Introduction

Observations of fast suprathermal electrons, due to the interaction of powerful radio waves with the ionosphere, using incoherent scatter radar (ISR) technique were first proposed and realized by Carlson, Wickwar and Mantas in 1972 (see Carlson et al. [1982]). The powerful O-mode wave in the resonance region near the reflection point effectively excites plasma waves – natural oscillations of the ionosphere plasma. Plasma waves due to the nonlinear cavitation process accelerate suprathermal electrons to energies in excess of $\varepsilon \sim 10\div 20$ eV, i.e. two orders of magnitude higher than the thermal electron energy $T_e \approx 0.1$ eV. The fast electrons propagate in the ionosphere to large distances, 10–100 km from the acceleration region. Those electrons excite nonthermal plasma waves, which were detected by ISR [Carlson et al., 1982].

We discuss here the results of two sets of experiments. The experimental data are here compared in detail with theory based on the mechanism of multiple electron acceleration.

Carlson et al. [1982] gave a qualitative explanation of their observations. They pointed out the significant role played by multiple scattering of electrons in the neutral atmosphere. The quantitative theory of multiple acceleration was proposed by Gurevich et al. [1985]. The comparison of the theory and observations can help reach a still higher level in our understanding of a number of space plasma phenomena in which accelerated electrons play an important role.

1.2 Observations

The most sensitive and direct ground based means of detecting 10–20 eV suprathermal electrons is by incoherent scatter plasma line observations of the Langmuir waves produced in the ionosphere by these electrons [Perkins and Salpeter, 1965]. This technique has been used to study photoelectron fluxes in the ionosphere [Yngvesson and Perkins, 1968]. Carlson et al. [1982] applied the ISR technique to measurements of suprathermal electrons produced by the interaction of powerful radio wave with the ionosphere at Arecibo. We present here two types of experimental data. The primary and secondary data were gathered on the nights 20 May 1972 and 13 July 1992 respectively. The primary data were published by Carlson et al. [1982], here we will revisit those results and supplement them with the complimentary set of secondary 1992 data [Djuth et al., 1996].

For the 1972 Arecibo data here, a HF dipole feed hung over the 1000-foot diameter ISR radar dishes, for which the 138 kW transmitter at 7.63 MHz delivered about 30 W/m² to the HF reflection height. For the 1992 Arecibo data here, a separate 48 array of log periodic antennas 17 km away from the ISR radar dish, for which the 300 kW transmitters at 5.1 MHz, into 50% efficiency transmission lines, delivered about 40 W/m² to the HF reflection height.

For 20 May 1972, strong plasma waves were excited near the reflection point at the altitude of $z = 285$ km. Night time plasma line intensities were observed to be enhanced by a factor of 10–100, over a range of altitudes which extended to below 250 km. They relaxed to their normal level within the electron transport time after the transmitter was turned off.

The 430 MHz incoherent scatter radar was used to diagnose the ionosphere, which allowed measurement of altitude profiles of the background plasma temperature T_e and electron concentration N . The plasma line echo from the plasma waves excited by the suprathermal electrons provided the altitude of a set of plasma frequencies f_p between 5 and 7.5 MHz, with an altitude resolution of 1.5 km.

For a given diagnostic radar wavelength λ_r the main scattered signal comes from ionosphere plasma waves, whose wave vector is directed toward the radar, and the phase velocity is equal to the velocity of suprathermal electrons v

$$v_{ph} = \frac{1}{2} \lambda_r f_p \quad (1)$$

Here $f_p = (e^2 N / 4\pi m)^{1/2}$ is the local plasma frequency. For the Arecibo radar $\lambda_r = 70$ cm and from eq. (1) one can obtain that the energy of electrons is

$$\varepsilon = \frac{1}{2} m v_{ph}^2 = 0.35 f_p^2 V \quad (2)$$

The intensities of plasma line echoes were shown in Fig.1 of Carlson et al. [1982]. A significant enhancement of radar echo intensities is evident at large distance below the reflection point of the radio wave at 285 km. The upgoing one-dimensional fluxes of suprathermal electrons at the heights 266 km and 256 km determined from these data are shown in Fig.3 of Carlson et al. [1982].

The additional data were gathered on the night 13 July 1992. The reflection altitude was 295 km, with strong enhancements of plasma line intensity observed in the vicinity of the reflection level. An example of plasma line echoes collected at the heights 337–367 km, where the ISR beam intersected the flux of the accelerated electrons, is shown in Fig.1.1. The echoes were obtained both below and above the maximum of F-layer. Electron density $N_{max} \approx 5.5 \cdot 10^5 \text{ cm}^{-3}$ was reached at F-maximum heights $z = 350$ km. The observed plasma line echoes correspond to upgoing electron flux. Note particularly here a significant enhancement of plasma line echoes even at altitudes more than 70 km above the reflection layer.

1.3 Brief Outline of the Theory

In ionosphere modification experiments, electrons gain energy in a strongly disturbed Langmuir resonance layer near the reflection point of the O-mode wave. The essential point of multiple acceleration theory is a large fraction of fast electrons, after leaving the acceleration layer will return back due to collisions with neutral particles, and thus can gain additional energy. The process can be repeated many times. Thereby, a wide region around the acceleration layer, which is elongated along the geomagnetic field, can be filled with strongly heated magnetized suprathermal electrons. The elongation along the

Earth's magnetic field can reach an order of a hundred kilometers because it is determined by two factors: the large mean free path of the suprathermal electrons and a small part of electron energy is lost in one collision. The extent of upward elongation will be considerably greater than in the downward direction, because the mean free path above the reflection height is considerably greater than below it. Consequently, a significant flux of suprathermal electrons moves to the magnetosphere, and can even reach the magnetic conjugate point.

The acceleration process due to the multiple crossing of acceleration layer is averaged and in final form, depends on two scalar factors only: the full power density P absorbed by fast electrons in the acceleration layer, and a characteristic parameter describing the effectiveness of the acceleration inside the layer T_{ef} . They are related to the details of the accelerating process, effective number of cavities, their width, and so on.

The distribution function of suprathermal electrons in the theory of multiple electron acceleration can be presented in the simple form [Gurevich et al., 1985]

$$f_0(\varepsilon, z) = CK_0\left(\frac{\varepsilon}{T_{ef}}\right) \exp\left\{-\left|\int_0^z \frac{dz}{L_e^\pm \cos \alpha}\right|\right\} \quad (3)$$

Where K_0 is a modified Bessel function, ε is the electron energy and T_{ef} is the effective temperature of suprathermal electrons, while α is the angle between the vertical and geomagnetic field. The normalization constant C is directly proportional to the power density P of the HF wave absorbed by the suprathermal electrons:

$$C = \frac{m^2}{4\pi^3 T_{ef}^3 \left(\frac{\delta}{3}\right)^{1/2}} P \quad (4)$$

Here δ is the average fraction of electron energy lost in a single collision with neutral molecules.

As given in eq. (3) the acceleration layer is assumed to be located at $z = 0$, while the L_e^\pm factor is the characteristic relaxation length of suprathermal electrons in the upward (+) and downward directions (-)

$$L_e^\pm(z) = \left[N_m^\pm \sigma_{tr}(\varepsilon) (3\delta)^{1/2} \right]^{-1} \quad (5)$$

Where $N_m^\pm = N_m^\pm(z)$ is the neutral density above (+) and below (-) the layer, $\sigma_{tr}(\varepsilon)$ is the total transport cross-section of electron-neutral collisions, and $\delta = \sigma_{in} / \sigma_{tr}$, where σ_{in} is the total cross-section of inelastic collisions, which includes ionization by electron impact. In σ_{tr} and σ_{in} the collisions with all neutral components are considered.

1.4 Comparison of the Observations with Theory

In the experiment on 20 May 1972, upgoing (in the radar look direction) fluxes of suprathermal electrons in the energy range $10 \div 17$ eV at the heights $z_1 = 256$ km and $z_2 = 266$ km were determined. The reflection point was at $z_0 = 285$ km. Comparison with the theory requires determination of the electron flux along radar direction e_r :

$$J_r = (\mathbf{v}, \mathbf{e}_r) f(\mathbf{v})$$

We introduce net upgoing and downgoing fluxes by integrating the flux J_r over corresponding angles in velocity space. Taking the angle between the geomagnetic field and vertical radar ray in Arecibo as $\alpha = 40^\circ$ we obtain the upgoing J_+ and downgoing J_- fluxes in energy interval $d\varepsilon$:

$$J_{\pm} = \frac{2\pi}{m^2} \mathcal{E}_0(\varepsilon) \left(1 \pm \frac{2}{3} (3\delta)^{1/2} \cos \alpha \right) \exp \left(- \left| \int_0^z \frac{dz}{L_{\varepsilon}^{\pm}(z) \cos \alpha} \right| \right) \quad (6)$$

The results of calculations are presented in Fig.1.2 for the given heights and different T_{ef} . It is apparent from the figure, that in the energy range 10 ÷ 20 eV reveal a rather flat spectrum. Besides, there is no strong difference between the fluxes J at different temperatures T_{ef} .

On the other hand, as follows from calculations the fluxes are effectively diminishing with distance from the acceleration layer. This is because electron energy is lost in inelastic collisions with the neutral. Equation (7) implies that the fluxes of upgoing electrons are smaller than the downgoing fluxes below the acceleration region. Their relation depends upon the angle α between radar ray direction and the geomagnetic field – for Arecibo at $\alpha = 40^\circ$ $J_+/J_- \approx 1.5$. Comparison between the theory and observations presented in Fig.1.3 shows a reasonable agreement between those two. In fact, the behavior of the spectrum of the electron flux at different heights is consistent with the theory for any of T_{ef} applied. Taking into account the absolute values of upgoing flux $J_+ = (4 \div 8) \times 10^5$ el/cm²s eV at characteristic energies 10 ÷ 15 eV obtained by Carlson et al. [1982], one can find the absorbed power W_S of the HF wave converted into acceleration of the suprathermal electrons

for	$T_{ef} = 5 \text{ eV}$	$W_S \approx 5.6 \text{ kW}$
for	$T_{ef} = 7.5 \text{ eV}$	$W_S \approx 6.2 \text{ kW}$
for	$T_{ef} = 10 \text{ eV}$	$W_S \approx 8.5 \text{ kW}$

In these calculations the Arecibo heater beam was taken roughly 10^3 km^2 for a full absorbed power $W_S = PS$. The factor P was determined from calculations of the electron flux made at different distances from the source, and from the results presented in Fig.1.3. It is apparent that the dependence of W_S on T_{ef} is not very strong. Note that eq. (7) allows us to estimate the upgoing electron flux at $z > 500 \text{ km}$ which escapes to the magnetosphere due to the absence of scattering collisions.

Figure 1.4 shows the height distribution of plasma line intensity measured during the second experiment, which is compared with the height dependence of the distribution function of suprathermal electrons $f_0(\varepsilon(z), z)$. Here energy $\varepsilon(z)$ is determined through the measured electron density distribution (Fig.1.1) by using eq. (2). One can see a sufficient agreement between the theory and observations. Note, that the normalized height dependence of the distribution function practically does not depend on the parameter T_{ef} (see Fig.1.4).

1.5 Discussion and Conclusions

It was shown that the ISR plasma line measurements are in reasonable agreement with the theory of multiple electron acceleration. From the observational data and their comparison with the theory it follows that: In experiments where powerful radio waves interact with the ionosphere a large number of suprathermal electrons are generated in the energy range up to 20 eV.

The suprathermal electrons are observed over a wide altitude region of the order of a few tens km both below and above the acceleration layer, centered near the reflection altitude of the powerful radio wave. While not observed, they must extend well above this altitude, into the plasmasphere/magnetosphere, and even to some degree into the magnetically conjugate region.

In power, $W_S \approx 5.6 \div 8.5 \text{ kW}$ (4 ÷ 6)% of the full transmitted heater power goes into the acceleration of suprathermal electrons. This energy is dissipated due to the

generation of optical emissions, ionization and heating of ionized and neutral components of the ionosphere plasma. The characteristic dissipation length for suprathermal electrons depends strongly on the altitude. For $z = 250$ km it is about 10 km, while for $z \approx 300$ km it is about 30 km.

The flux of suprathermal electrons into the plasmasphere/magnetosphere depends strongly on the height z_0 of the acceleration layer. For the studied cases $z_0 = 285$ km, and $z_0 = 295$ km, and $W_S = 8.5$ kW, $T_{ef} = 10$ eV the flux into magnetosphere in the energy range $10 \div 20$ eV is about 0.9 kW. This estimate is consistent with earlier calculations of photoelectron energy and flux escape into and loss to the plasmasphere/magnetosphere [Mantas et al., 1978].

We conclude that plasma line observations combined with the theory enables us to obtain significant information about the acceleration of the suprathermal electrons in ionosphere modification experiments. For the first time the structure and size of the perturbed region filled with suprathermal electrons along with the full power going into the accelerated electrons is determined by a theoretical calculation consistent with an observational data set.

Note that we had the possibility to use here only a small part of ISR data on ionosphere existing from experiment 1992. The further elaboration of this data and their comparison with the theory is of a significant interest and will be definitely quite fruitful.

We suggest also these experiments be repeated to obtain more information about the main features of acceleration, and its dependence upon the reflection height and ionosphere conditions. Note that such observations have as yet been done only at Arecibo, though we expect that the latitude dependence of the effect could be very significant.

At the same time, the agreement that found between the theory and observations make it possible to plan more detailed future experiments. These will thus help to reach a much better understanding of the physical mechanisms of electron acceleration and nonlinear processes in ionosphere plasmas.

1.6 References

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Figure Captions

Figure 1.1 Spectra of the downshifted plasma line measured at Arecibo 13 July 1992 at 8:30 LT.

Figure 1.2. Upgoing flux of suprathermal electrons versus their energy computed for the full absorbed power $W_S = 10$ kW and for different values of T_{ef}

Figure 1.3. Upgoing flux of suprathermal electrons versus their energy. Solid curves 1, 2 and 3 correspond to $T_{ef} = 5, 7.5$ and 10 eV respectively, while the points with bars correspond to the observations made on 052072.

Figure 1.4. Plasma line intensities versus altitude. Solid curves 1, 2 and 3 correspond to $T_{ef} = 5, 7.5$ and 10 eV respectively, while the points with bars correspond to the observations made on 071392.

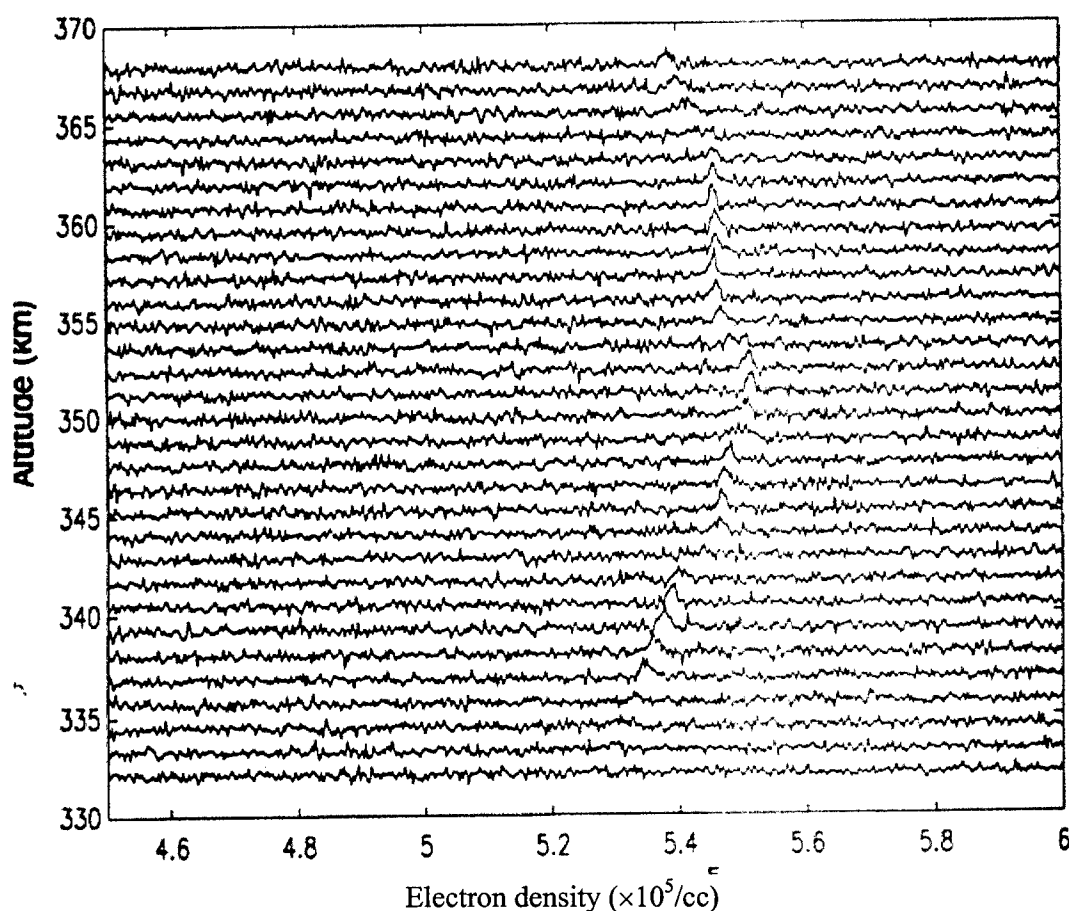


Figure 1.1

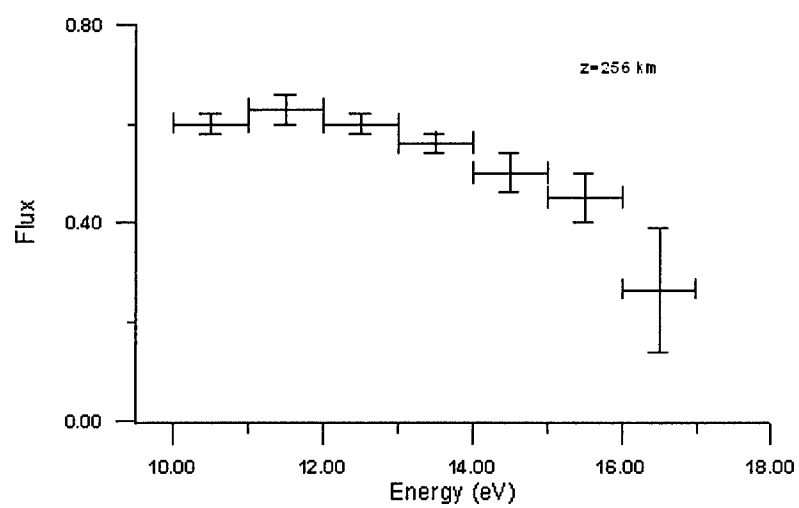
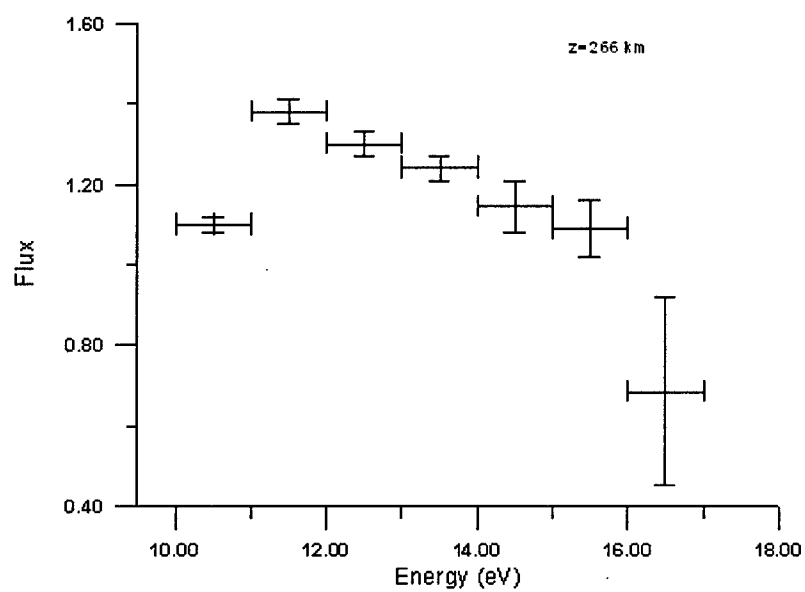


Fig.1.2

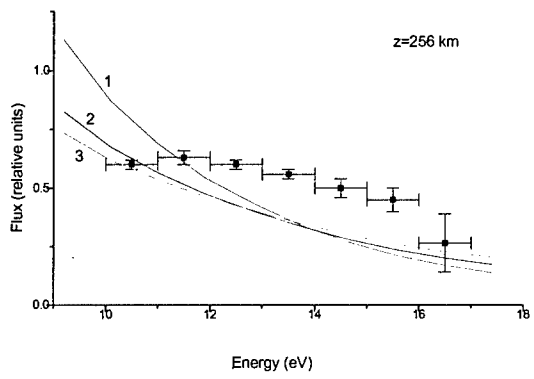
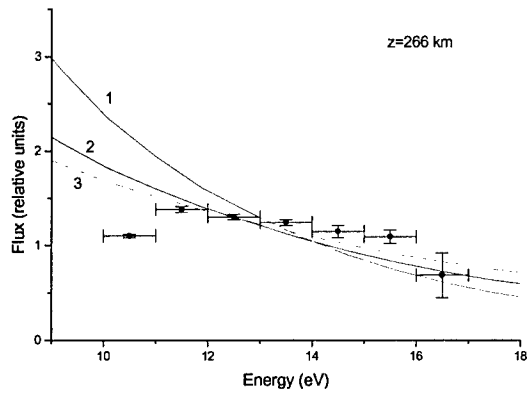


Figure 1.3

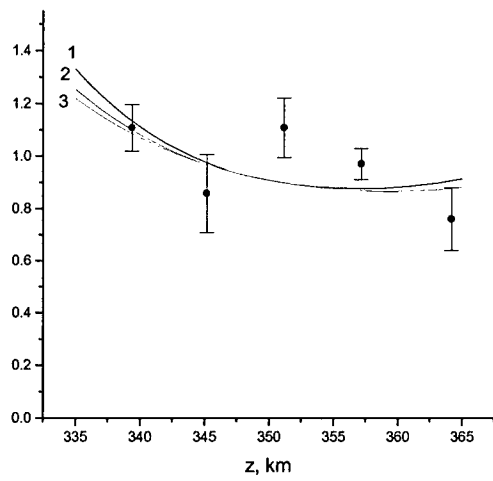


Figure 1.4

2. Optic emission from modified by powerful radio waves ionosphere

Introduction

Enhancement of optic emission was established in the first ionosphere modification experiments (Sipler and Biondi, 1971; Adeishvili et al., 1976). It was studied afterwards in multiple observations (Haslett and Megill, 1974; Carlson, 1974; Bernhardt et al., 1989). During the last years new interesting results have been obtained at HAARP facility (Peterson, 1999; Carlson, 1999), in sporadic E-layer at Arecibo and at low duty circle experiments at SURA (Nazyrov et al., 1999). The optic emission measurements indicate directly the existence of suprathermal electrons accelerated up to the energies 10 eV and more in modified ionosphere.

Other type of observations, demonstrating the accelerated up to the energies 20 eV electrons was performed by Carlson et al., 1982, Juth et al., 1996, using incoherent scattering radar at Arecibo. The suprathermal electrons were seen in these experiments in a wide region of the order 100-km around the reflection point of the heater wave. Some indications of the significant effect of suprathermal electrons on SEE emission exist as well (Frolov, 1998).

Theoretical considerations from the very beginning connected the enhancement of optic emission with the acceleration of electrons in the Langmuir plasma turbulence layer near the reflection point of powerful HF O-wave. Though, it was indicated as well, that the mostly intensive red-line emission (630 nm) partly could be determined by the direct effective heating of the main bulk of electrons (Carlson, Mantas, 1996), or it's significant part, captured inside nonlinear density depletions – "striations" generated by the upper-hybrid turbulence (Gurevich and Milikh, 1998). The theory of Langmuir plasma turbulence generated near the reflection region of pump O-wave connects the electron acceleration process with the modulation instability and with the creation of nonlinear Langmuir cavitons – density depletions filled with Lagmuir plasma oscillations (Wong and Stenzel, 1975). The fast electrons, passing trough the turbulent region filled with cavitons obtain the energy from the trapped caviton plasma oscillations in a diffusive way (Eastebrook et al., 1975; Morales and Lee, 1977; Andreev et al., 1980; Wang, Goldman and Newman, 1997). The main feature of this mechanism is that the energy is gained only by fast electrons whose velocity is high enough $v > a/\omega$, where a is a scale of a caviton, ω is the heater wave plasma frequency. Low energy electrons oscillate in the caviton adiabatically and do not get any additional energy. It should be emphasized, that the collapse of cavitons does not effect strongly the acceleration process of fast electrons (Wang et al., 1997). In the same time collapse can input a significant part of wave energy to the thermal component (Du Bois et al. 1993, Hanssen et al. 1997). To calculate the flux of energy, which goes to fast electrons the detailed, kinetic theory is needed, which takes into account electron collisions.

For example, the number and energy distribution of fast electrons in ionosphere conditions is affected by the so-called "multiple acceleration" (Gurevich et al., 1985, Vaskov et al., 1983). The multiple acceleration process is determined by the fact, that the Langmuir acceleration take place in a strongly excited layer, embedded in a weakly disturbed ionosphere plasma. After the fast electrons leave the acceleration layer it collides with a neutral molecules in nondistrubed plasma. Due to collisions some electrons could return to the acceleration layer and get there the energy once more. The process could be repeated many times, thus creating a "multiple acceleration". Directly this process determines the high-energy tail of electron distribution function and a wide spread of fast electrons around the acceleration region.

Though, understanding of the main mechanisms of the acceleration process in Langmuir turbulence seems significantly clear enough, both the theory and observations in ionosphere modification experiments up to now were not developed to the state when a detailed quantitative comparison between them could take place. That means, that some fundamental problems of modification experiments remain non-answered. One of them is the intensity and space distribution of the observed optic emission.

In the present work we intend to develop the existing theory and apply it to obtain the detailed comparison with the main part of the existing experimental data. The paper is constructed as follows. In section 2 heating of electrons by both in upper hybrid and in Langmuir turbulence is considered. It is demonstrated that the calculations of electron acceleration in Langmuir resonance layer without electron collisions lead to controversial result. In section 3 the theory of multiple acceleration is generalized. The energy losses of electrons on the excitation of molecular vibration level are taken into account. In section 4 the model calculations describing the general peculiarities of the suprathermal electron space energy distribution are performed. The section 5 describes the detailed comparison of the theory with the main observational data of the artificial optic emission in F-layer. The strong enhancement of the optic emission in sporadic E-layer modification is discussed in sections 6. In conclusion we formulate the nowadays state of the problem. The strongly inhomogeneous character of the disturbed region is stressed.

1.2 Heating of electrons in the resonance layer

In the vicinity of reflection point of powerful O-wave in ionosphere both heating and acceleration of plasma electrons takes place.

As is well known for vertically propagating pump O-wave the resonance region is situated at the heights Z between upper hybrid Z_{uh} and Langmuir Z_L resonances

$$Z_{uh} \leq Z \leq Z_L \quad (1)$$

In this region the frequency of the pump wave ω coincide with the frequency of the natural oscillations in ionosphere plasma

$$\omega = \omega_L = \sqrt{\frac{4\pi e^2 N(Z_L)}{m}} \quad (2)$$

Here $N(Z_L)$ – electron density at the Langmuir resonance, where the Langmuir turbulence take place. And Z_{uh} is defined by resonance condition for upper hybrid resonance

$$\omega = \omega_{ph} = \sqrt{\omega_L^2 + \omega_c^2} \quad (3)$$

In the vicinity of Z_{uh} the upper hybrid (UH) turbulence is effectively developing. The width of resonance region ΔZ depends on the pump wave frequency and plasma density gradient L :

$$\Delta Z = Z_L - Z_{uh} = L \frac{\omega_c^2}{2\omega^2} \quad L = \frac{N}{|dN/dZ|} \quad (4)$$

ω_c is electron gyrofrequency. For the ionosphere F-layer conditions usually $\Delta Z \sim 2 \div 5$ km.

Conventional collision absorption of radio waves in F-region is very weak – less than 1 dB [Ginzburg, 1967]. If plasma turbulence in the resonance region is excited the absorption become abnormally strong – up to 10÷20 dB [Gurevich, 1978].

Note, that the Langmuir turbulent layer (LT) is placed in the upper part of resonance region close to the reflection point (2), while close to the upper hybrid resonance (3) lays upper hybrid turbulent layer (LH). Acceleration of electrons take place

in Langmuir resonance layer (we do not speak here about singled out case of multiple gyroresonance frequency $\omega = n\omega_c$). The UH turbulence lead to formation of striations. An effective anomalous absorption of the pump wave take place in this region, which lies below the Langmuir turbulence layer. It means that the UH anomalous absorption (absorption on striations) can affect significantly the heating and acceleration processes in the Langmuir turbulence layer while diminishing the pump-wave power.

Heating of electrons in UH turbulence region

Anomalous absorption is determined by the excitation of UH waves due to direct transformation of the pump wave on the density gradient in striations. UH waves trapped in striations heat electrons and thus determine the depth form and number density of striations. Nonlinear theory of the electron heating and anomalous absorption on striations was developed by Gurevich et al. 1996. The results of the theory are in a rather good agreement with observations.

Following to the theory of UH anomalous absorption we determined the simple relation for the losses of pump wave power P in the UH turbulence region on its way to Langmuir turbulent layer.

$$P(Z_L) \approx P_0 10^{-K/10} \quad K = \gamma \left(\frac{L}{100 \text{ km}} \right) Q$$

$$\gamma \approx 5.5 \div 8, \quad 0 \leq Q \leq 1$$

Here K is an attenuation coefficient in dB. Coefficient Q is proportional to the relative depth of striations. It depends on pump wave power, frequency, plasma parameters. The state of the striation development is also significant. For the well developed striations $Q \sim 1$.

1.3 Heating of electrons in Langmuir turbulent layer

The excitation of Langmuir turbulence, its steady state and caviton formation in the field of powerful radio wave in conditions close to ionosphere was studied numerically. We will discuss the result of the calculations of Hanssen et al and Wang et al in order to determine the thermal heating of electron by Langmuir turbulence. The parameters used in calculations are presented in Table 1.

	Hanssen et al (1992)	Wang et al (1997)
Pump wave frequency (MHz)	f= 4.01	f=4.9
Pump wave field E_{rms} (V/m)	1.26	1.6
Plasma density n_0 (cm^{-3})	2×10^5	2×10^5
Electron temperature (eV)	0.17	0.1
Pump wave energy $W_{rms} = E_{rms}^2 / 4\pi n_0 T_0$	2.7×10^{-3}	5×10^{-3}

Hanssen et al. (1992) considered one-dimensional Zakharov system of equations describing nonlinear development of plasma density and electric field fluctuations under the action of homogeneously oscillating electric field of the pump wave. Electron and ion collisions and Landau damping were taken phenomenologically, using terms from linear theory. The established average energy of longitudinal plasma waves, concentrated mostly in cavitons grew up 70 times in comparison with the pump wave energy. It gives $W_{rms}^p \approx 0.19$.

The power density dissipating by plasma oscillations to the heating of electrons is

$$\frac{dW}{ds} = \frac{e^2 E_{rms}^2 v_e n_0}{m \omega^2} \quad (5)$$

Where s is coordinate along the Earth magnetic field, v_e – electron collision frequency, ω – plasma wave frequency and $E_{rms}^2 = 4\pi n_0 T_e W_{rms}$. Following Hanssen et al. let us suppose that the whole width of the turbulent layer is of the order of pump wave length near reflecting point $L \sim \lambda \approx 200$ m. We obtain that the full pump wave power dissipating in Langmuir cavitons to the heating of plasma electrons:

$$W_T \approx W_{rms} v_e T_e \lambda_m \approx 40 \frac{\mu kW}{m^2} \quad (6)$$

That value is quite significant. Really, if we neglect anomalous absorption in striations below turbulent layer, then to create the electric field $E_{rms} = 1.26$ V/m (see Table 1) at the main Eiry maximum near reflection point, the pump wave power $60(\sin\alpha)^{4/3} (100\text{km}/L)^{1/3} \mu kW/m^2$ is needed. Here α is the inclination angle of Earth magnetic field. We see that at $\alpha \sim 1$ about 5 dB of pump wave is absorbed due to the thermal heating of electrons in the cavitons. Note that the calculated pump wave absorption is effectively growing with latitude ($\alpha^{-4/3}$). So, more than 60% of absorbed pump-wave energy in caviton turbulence is going directly to the thermal component of electrons.

Wang et al. 1997 elaborated one-dimensional kinetic model. Not only field equation but also Vlasov kinetic equation for noncolliding electrons was integrated simultaneously in the model. The Langmuir waves instability was excited and established. Energy level of average plasma oscillations grew up about 30 times in comparison with pump wave field reaching $W_{rms} = 0.14$. We see that in spite of quite different approach, the numerical value of saturated energy of plasma oscillations is close enough in both calculations.

The power density dissipating to the thermal component of electrons is again determined by expression (6). The width of turbulent layer was supposed to be $\lambda_m \approx 180$ m. From (6) we again obtain the heating power $W_T = 40 \mu kW/m^2$, what means that about 2.5 dB of pump wave energy goes to the heating in cavitons.

We see that when Langmuir turbulence is effectively excited the pump wave energy losses to the heating of electrons is quite significant, compatible (though remaining less) to the anomalous absorption in UH layer.

Significant advantage of Wang et al. 1997 kinetic model is the possibility to see directly not only saturation of plasma oscillations, but acceleration of suprathermal electrons as well. The result of calculations shows that a tail of initial Maxwellian distribution function is growing in time. An established suprathermal stationary tail is

$$f_s = 8.3 \times 10^{-6} \left(\frac{32V_{Te}}{v} \right)^{1.9} \frac{n_0}{V_{Te}} \quad 4 \leq \frac{v}{V_{Te}} \leq 25 \quad (7)$$

shown to be approximated by power – law:

Here $V_{Te} = (T_e/m)^{1/2}$ – thermal velocity. The one dimensional maxwellian part of distribution function has standard view:

$$f_m = \frac{n_0}{\sqrt{2\pi} V_{Te}} \exp\left(-\frac{v^2}{2V_{Te}^2}\right) \quad (8)$$

The full number of suprathermal electrons is not large about 2% of the thermal electrons.

Let us determine now the transport of energy by fast electrons. Integrating distribution function (7) we can calculate the full power transported out of the turbulent layer by accelerated suprathermal electrons

$$W_a = 2 \int_{4V_{Te}}^{25V_{Te}} \frac{1}{2} m v^2 v f_s dv \approx 2.68 n_0 T_e V_{Te} \approx 1700 \frac{\mu kW}{m^2} \quad (9)$$

We see that the power transported by suprathermal electrons is 20 times larger, than the full power of the pump wave ($W_p \approx 96 \mu kW/m^2$). This result is obviously controversial. It means, that in spite of the fact that mean free path of the suprathermal electrons in ionosphere F-layer is larger than the width of the turbulent region $\lambda_m \approx 180$ m, the electron collisions could not be neglected. The flux of the energy out of the layer is determined by collisions – suprathermal electrons not only leave the layer, but due to collisions they can return back. Thus only the full balance between accelerated inside the layer and returning backward suprathermal electrons can determine correct relation between the pump wave energy losses and acceleration. Just this balance is studied in the theory of multiple acceleration. This theory in modernized form would be considered below.

Note, that the controversial result of electron acceleration in non collisional kinetics shows, that the theory of plasma turbulence saturation should be revised as well: electron collisions outside the Langmuir turbulent layer should be taken into account to organize a correct flow of plasma energy. In other words boundary conditions, which connect turbulent and nonturbulent regions, could affect the process of Langmuir waves saturation inside the turbulent layer.

1.4 Kinetic theory of multiple acceleration

To find the distribution function we have to solve kinetic equation inside and outside acceleration layer and match solutions.

Inside the layer kinetic equation has a form:

$$\frac{\partial f}{\partial t} + v_{\parallel} \frac{\partial f}{\partial z} - \frac{e}{m} E \frac{\partial f}{\partial v_{\parallel}} = 0 \quad (10)$$

Here E is electric field of Langmuir cavitons excited by powerful radio wave

$$E(z, t) = \frac{1}{2} E(z) (e^{-i\omega t} + c.c) \quad (11)$$

The solution of the equation (10) has the form:

$$f(z, \varepsilon_{\parallel}, t) = F(z - Z(t), \varepsilon_{\parallel} - \varepsilon(t)) \quad (12)$$

Here $\varepsilon(t)$, $Z(t)$ the trajectory of a fast particle in electric field E :

$$\frac{dZ}{dt} = V, \quad \frac{dV}{dt} = \frac{e}{m} E, \quad V = \sqrt{2\varepsilon/m}$$

To solve our problem we need to determine the changing of distribution function Δf when a fast particle crossing the acceleration layer. Averaging relation (12) on fast oscillations we find the changing of distribution function after passing through the acceleration layer).

$$\Delta f = \frac{1}{2} \frac{\partial}{\partial \varepsilon_{\parallel}} \left[\langle (\Delta \varepsilon)^2 \rangle \frac{\partial F}{\partial \varepsilon_{\parallel}} \right] \quad (13)$$

According to (11) after averaging on fast oscillations in a fix spatial point energy gain is given by a formulae:

$$\Delta \varepsilon(z, t) = -\Delta \varepsilon \cos \left(\omega t - \frac{\omega z}{v_{\parallel}} + \psi \right) \quad (14)$$

$$\Delta \varepsilon(\varepsilon_{\parallel}) = e|E_k|, \quad E_k = \int_{-\infty}^{+\infty} E(z) e^{-ikz} dz, \quad k = \frac{\omega}{v_{\parallel}} = \omega \sqrt{\frac{m}{2\varepsilon_{\parallel}}}$$

For example, for Gauss like caviton $E(z) = E_0 \exp(-(z/a)^2)$ we have from (11):

$$\Delta \varepsilon = \pi^{1/2} e E_0 a \exp(-ka^2) \quad (15)$$

We see from (15) that only electrons with high enough energy

$$\varepsilon_{\parallel} > \varepsilon_1 = m(\omega a)^2 = T_e \left(\frac{a}{D_e} \right)^2 \quad \left(\frac{a}{D_e} \right)^2 \gg 1 \quad (16)$$

where D_e is a Debay length, could be accelerated. It is easy to see that relation (16) has a universal character and does not depends on a real form of a caviton. It is obvious that if we have some number of noncorrelated cavitons the average energy changing is equal to a sum

$$\langle \Delta \varepsilon^2 \rangle = \sum_i \langle \Delta \varepsilon_i^2 \rangle$$

Now let us consider the solution outside acceleration layer at $z > 0$ and $z < 0$. Kinetic equation has a form:

$$\frac{\partial f}{\partial t} + v \mu \frac{\partial f}{\partial z} = -S \quad (17)$$

Here $f(t, v, \mu, z)$ is the distribution function of electrons, S - collision integral, μ - cosine of an angle between the velocity of electron and z direction. The collision integral consists of two parts: inelastic S_0 and elastic S_1 . The collision integral depends strongly on energy. For energies $\varepsilon < \varepsilon^*$ ($\varepsilon^* \approx 4$ eV) the main collision crosssection is defined by oscillational levels and inelastic collision integral has a diffusive form. But for energies $\varepsilon > \varepsilon^*$ the collision crosssection is determined by optical levels. Thus the collision integral could be presented in a form:

$$S_0(F_0) = -\frac{\partial}{\partial \varepsilon} \left(R \frac{\partial f}{\partial \varepsilon} \right) \quad \varepsilon < \varepsilon^* \quad (18)$$

$$S_0(F_0) = \nu_0(\varepsilon) F_0 \quad \varepsilon > \varepsilon^* \quad \nu_0 = \nu \sigma_0 N_m$$

$$S_1(f_1) = \nu_1(\varepsilon) f_1, \quad \nu_0 = \nu \tilde{\sigma}_i N_m$$

$$\tilde{\sigma}_i = \sigma_i \theta(\varepsilon - \varepsilon^*) + \sum_i \theta(\varepsilon - \varepsilon^*)$$

The σ_0 is the sum of effective crosssection of inelastic collisions with neutral atoms N_m - the total density of neutrals. σ_i is transport electron crosssection of electron elastic collisions

$$\sigma_0 = \sum_k \sigma_{k0} \frac{N_{mk}}{N_m} \quad \sigma_i = \sum_k \sigma_{ik} \frac{N_{mk}}{N_m} \quad N_m = \sum_k N_{mk} \quad (19)$$

The averaged part of total energy of an electron loosing in one collision is small; it is defined by parameter

$$\delta = \nu_0 / \nu_1 = \sigma_0 / \sigma_i \quad (20)$$

High electron energies ($\varepsilon > \varepsilon^*$)

Because of (20) we can expand the equation (17) on small parameter δ . At first approximation the distribution function does not depend on μ . Taking into account first approximation one can find:

$$\frac{\partial F_0}{\partial t} + \frac{v}{3} \frac{\partial f_1}{\partial z} = -S_0, \quad S_0 = \frac{1}{2} \int_{-1}^1 S d\mu \quad (21)$$

$$\frac{\partial F_0}{\partial z} = S_1 = -v_1 f_1$$

Boring in mined the boundary conditions (13) at $z = 0$ in stationary conditions from (21) one can find:

$$\frac{\partial F_0}{\partial z} = -\text{sign}(z) \frac{3v_1}{9v\varepsilon} \frac{\partial}{\partial \varepsilon} \left\{ \int_0^\varepsilon \langle \Delta \varepsilon^2(\varepsilon_{\parallel}) \rangle d\varepsilon_{\parallel} \frac{\partial f_0}{\partial \varepsilon} \right\} \quad (22)$$

It is easy to see that for collision integrals in the form (18) the equation (22) depends on energy parametrically only. So we can rewrite equation (18) in the region of energies

$$\frac{\partial^2 F_0}{\partial \xi^2} = 3\delta F_0, \quad \xi = \int_0^\varepsilon \frac{v}{v_1} dz = \int_0^\varepsilon \sigma_1 N_m(z) dz \quad (23)$$

Taking into account that δ is a slow changing function it is possible to solve (23) in WKB approximation

$$F_0(\varepsilon, z) = F_0(\varepsilon) \left[\frac{\delta(0)}{\delta(z)} \right]^{1/4} \exp \left\{ - \left| \int_0^z \frac{dz}{L_\varepsilon(z)} \right| \right\} \quad (24)$$

$$f_1(\varepsilon, z) = \text{sign}(z) (3\delta)^{1/2} F_0(\varepsilon, z)$$

Here L_ε is a relaxation scale of fast electrons with given energy

$$L_\varepsilon = \frac{v}{(3v_1 v_0)^{1/2}} = \frac{1}{N_m (3\sigma_1 \sigma_0)^{1/2}}$$

Middle electron energies ($T_e < \varepsilon < \varepsilon^*$)

In the region of energies $\varepsilon < \varepsilon^*$ the kinetic equation (21) takes the form:

$$\frac{\partial^2 F_0}{\partial \xi^2} = \frac{\partial}{\partial \varepsilon} \left(R \frac{\partial F_0}{\partial \varepsilon} \right) \quad (25)$$

Since the solution of kinetic equation in the region $\varepsilon > \varepsilon^*$ has a form (24), to match this solution with solution in region $\varepsilon < \varepsilon^*$ we had to search the solution for $\varepsilon < \varepsilon^*$ in a form

$$F_0 = g(\varepsilon) \exp \left\{ - \left| \int_0^z \frac{dz}{L_\varepsilon(z)} \right| \right\}$$

The equation (25) takes the form:

$$\frac{\partial}{\partial \varepsilon} \left(R \frac{\partial g}{\partial \varepsilon} \right) = - \frac{v_0 v_1}{\tilde{v}_1} g$$

Taking into account that

$$v_0 v_1 \ll R \tilde{v}_1$$

One can see from (26) that at first approximation

$$F_0(\varepsilon) = F_0(\varepsilon^*) = \text{const} \quad \varepsilon < \varepsilon^* \quad (27)$$

The distribution function $F_0(\varepsilon) = F_0(\varepsilon, z)|_{z=0}$ according to boundary condition (13) is determined by diffusion equation:

$$\frac{1}{\delta^{1/2} \varepsilon} \frac{d}{d\varepsilon} \left\{ \delta^{1/2} \varepsilon T_{\text{eff}}^2 \frac{df_0}{d\varepsilon} \right\} = f_0 - f^{(0)}(\varepsilon), \quad F_0 = f_0 - f^{(0)}(\varepsilon) \quad (28)$$

$$\text{and } F_0(\varepsilon) = F_0(\varepsilon^*) = \text{const}, \quad \varepsilon < \varepsilon^*$$

We introduce here effective temperature :

$$T_{\text{eff}}(\varepsilon) = \kappa \bar{\Delta} \varepsilon, \quad \bar{\Delta} \varepsilon = \left\{ \frac{1}{\varepsilon} \int_0^\varepsilon \langle \Delta \varepsilon^2(\varepsilon_{\parallel}) \rangle d\varepsilon_{\parallel} \right\}^{1/2}, \quad \kappa = \frac{3^{1/4}}{2\sqrt{2}} \delta^{-1/4} \quad (29)$$

here $\bar{\Delta} \varepsilon$ is mean electron energy increasing in accelerating layer. The effective temperature increase gradually with increasing energy and saturate at the level.

$$T_{\text{effm}} = \kappa (n/2)^{1/2} \Delta \varepsilon_m, \quad \varepsilon_m \leq \varepsilon_1$$

Here n is the number of cavitons in accelerating layer.

The intensity of electron acceleration is defined by parameter

$$\gamma = \frac{T_{\text{eff}}}{T}, \quad T = -f^{(0)}(\varepsilon) \left/ \frac{df^{(0)}(\varepsilon)}{d\varepsilon} \right.$$

For weak acceleration ($\gamma < 1$) the distribution function of electrons does not change substantially

$$F_0(\varepsilon) = \left(\frac{T_{\text{effm}}}{T} \right)^2 f^{(0)}(\varepsilon)$$

In opposite case of strong acceleration $\gamma \gg 1$ the equation (29) in WKB approximation has the following solution:

$$f_0(\varepsilon) = f^{(0)}(\varepsilon) \quad \varepsilon < \varepsilon_{th}$$

$$f_0(\varepsilon) = f_0(\varepsilon_{th}) \left\{ \frac{\pi \varepsilon_{th} T_{\text{effm}} \delta^{1/2}(\varepsilon_{th})}{2 \varepsilon T_{\text{effm}}(\varepsilon) \delta^{1/2}(\varepsilon)} \right\}^{1/2} \left\{ \frac{dT_{\text{effm}}}{d\varepsilon} \right\}^{-1/2} \exp \left[- \int_{\varepsilon_{th}}^{\varepsilon} \frac{d\varepsilon}{T_{\text{effm}}(\varepsilon)} \right] \quad \varepsilon > \varepsilon_{th} \quad (30)$$

The value ε_{th} in (30) is defined by matching of distribution functions in saddle point

$$T_{\text{effm}} = T(\varepsilon_{th}).$$

In the case of strong acceleration in the limit $T_{\text{eff}} \approx \text{const}$, $\delta^{1/2} \approx \text{const}$ the distribution function (30) has the form:

$$f(\varepsilon) = CK_0 \left(\frac{\varepsilon}{T_{\text{eff}}} \right) \quad (31)$$

Here $K_0(x)$ is a modified Bessel function. The number of fast electrons is defined by the integral

$$N_e^f = \int f(\varepsilon) d^3v$$

But the normalized constant C is determined by the microwave power absorbed by fast electrons. The absorbed power consists of two terms

$$\begin{aligned}
P &= P_1 + P_2 & P_1 &= \frac{16\pi}{3m^2} \left(\delta/3 \right)^{1/2} C \varepsilon_{th}^3 K_0 \left(\frac{\varepsilon_{th}}{T_{eff}} \right) \\
P_2 &= \frac{16\pi}{m^2} \left(\delta/3 \right)^{1/3} C \int_{\varepsilon_{th}}^{\infty} \varepsilon^2 K_0 \left(\frac{\varepsilon_{th}}{T_{eff}} \right) d\varepsilon
\end{aligned} \tag{32}$$

The first term P_1 in (32) is responsible for losses on rotational levels of molecules and the second one P_2 describes the electron acceleration.

4. Drift effects on the scale of field-aligned irregularities

Irregularities of electron temperature and plasma density in the F-region of the ionosphere are strongly elongated magnetic field due to dramatic difference in transport coefficients along the magnetic field and transverse to it. Their scales along the magnetic field and reply to heating sources in stationary state were determined previously neglecting plasma drift effects both in linear [Gurevich, 1978] and many other publications] and nonlinear [Gurevich et al., 1995] approximations. These scales are rather long reaching several ten kilometers. Everybody settle that such elongated irregularities should be affected by ionospheric drifts perpendicular to magnetic field, but up to date there was no theoretical estimations for this influence. It can be just supposed that drifts could substantially diminish field-aligned scales and efficiency of excitation of irregularities with sufficiently small scales in the drift direction. Here we will attempt to consider drift influence on electron temperature and plasma density irregularities under approximation that source dimension along magnetic field is much less than characteristic transport lengths, and neglecting all transport processes across magnetic field. Two main problems will be considered: constant drift effects on irregularities caused by fixed in the ionosphere sources (for example, narrow radio beam), and gradient drift effects on irregularities created by moving with plasma sources (like excitation of irregularities by plain electromagnetic wave).

Constant drift velocity

Transport equations along the magnetic field for disturbances of electron temperature T and number density N can be written taking into account homogeneous drift with constant velocity V_0 across magnetic field in the form [Gurevich, 1978]

$$\kappa_{\parallel} \frac{\partial^2 T}{\partial z^2} - V_0 \frac{\partial T}{\partial x} - \delta \nu T = -Q_T \delta(z) \quad (1)$$

$$D_{\parallel} \frac{\partial^2 N}{\partial z^2} - V_0 \frac{\partial N}{\partial x} - \gamma N = - \left(Q_N - \frac{D_{T\parallel}}{\kappa_{\parallel}} Q_T \right) \delta(z) - \frac{D_{T\parallel}}{\kappa_{\parallel}} \left(V_0 \frac{\partial T}{\partial x} + \delta \nu T \right)$$

Here D_{\parallel} , $D_{T\parallel}$, and κ_{\parallel} are diffusion, thermal diffusion, and heat conductivity along the magnetic field respectively, γ is the recombination coefficient, δ is the average part of electron energy loss under collisions with ions and neutrals, and ν is the electron collision frequency. It is supposed in (1) that irregularities are induced by fixed heating Q_T and density Q_N sources which scales along the magnetic field is much less than characteristic scales of transport processes so that Dirac delta function $\delta(z)$ can be used. Coordinate z is along the magnetic field, and x is along drift velocity V_0 perpendicular to the magnetic field.

Introducing un-dimensional variables according to

$$\tau = \frac{T}{T_0}, \quad n = \frac{N}{N_0}, \quad \xi = \frac{x}{L_{\nu}}, \quad \varsigma = \frac{z}{L_T}, \quad (2)$$

with characteristic lengths

$$L_T = \sqrt{\frac{\kappa_{\parallel}}{\delta \nu}}, \quad L_{\nu} = \frac{V_0}{\delta \nu} \quad (3)$$

and un-dimensional heating and density sources as

$$q_T = \frac{1}{T_0 \sqrt{\kappa_{\parallel} \delta v}} Q_T, \quad q_N = \frac{1}{N_0 \sqrt{\kappa_{\parallel} \delta v}} \frac{\kappa_{\parallel}}{D_{\parallel}} Q_N, \quad (4)$$

where T_0 and N_0 are undisturbed electron temperature and number density respectively, and looking for a solution of (1) in the form $(\tau, n) \propto \exp(ik\xi)$, one can rewrite transport equations (1) as

$$\frac{d^2 \tau}{d\zeta^2} - (ik + 1)\tau = -q_T \delta(\zeta), \quad (5)$$

$$\frac{d^2 n}{d\zeta^2} - (i\beta k + \mu^2)n = -(q_N - Rq_T)\delta(\zeta) - R(ik + 1)\tau$$

Following denotations are used in (5):

$$\beta = \frac{\kappa_{\parallel}}{D_{\parallel}}, \quad \mu = \frac{L_T}{L_N} = \sqrt{\frac{\kappa_{\parallel} \gamma}{v D_{\parallel}}}, \quad R = \frac{T_0 D_{T\parallel}}{N_0 D_{\parallel}}. \quad (6)$$

Solution of (5) depends on $|\zeta|$ (that can be proved by substitution $\zeta \rightarrow -\zeta$), and only a region of $\zeta > 0$ will be considered. In this case (5) takes a form of

$$\frac{d^2 \tau}{d\zeta^2} - (ik + 1)\tau = 0, \quad (7)$$

$$\frac{d^2 n}{d\zeta^2} - (i\beta k + \mu^2)n = -R(ik + 1)\tau$$

with boundary conditions

$$\tau|_{\zeta \rightarrow \infty} = 0, \quad 2 \frac{d\tau}{d\zeta} \Big|_{\zeta \rightarrow 0} = -q_T, \quad (8)$$

$$n|_{\zeta \rightarrow \infty} = 0, \quad 2 \frac{dn}{d\zeta} \Big|_{\zeta \rightarrow 0} = -(q_N - Rq_T).$$

Solution of (7), (8) is straightforward and leads to

$$\tau = \frac{1}{2p} q_T e^{-p\zeta}, \quad n = \frac{1}{2s} q_T e^{-s\zeta} - \frac{1}{2(p^2 - s^2)} Rq_T (pe^{-p\zeta} - se^{-s\zeta}), \quad (9)$$

where $p = \sqrt{1 + ik}$ and $s = \sqrt{\mu^2 + i\beta k}$ and $\mathcal{R}(p), \mathcal{R}(s) > 0$ (\mathcal{R} stands for the real part).

Let us analyze (9) supposing for simplicity $q_N = 0$. Amplitudes of electron temperature and density irregularities produced by constant heating source q_T depend on its scale along drift direction decreasing for smaller scales. The same dependence exists for the length of irregularities along the magnetic field. There is also phase shift between heating source and irregularities. Temperature disturbances can be characterized by "efficiency" of heating G_T and their relative to heating sources phases Φ_T that can be defined as

$$G_T = \frac{2}{q_T} |\tau_0| = \frac{1}{|\sqrt{1 + ik}|} = (1 + k^2)^{-1/4}, \quad \Phi_T = \arg(\tau_0) = \arg \frac{1}{\sqrt{1 + ik}} = -\frac{1}{2} \arctan k. \quad (10)$$

Similar parameters of "efficiency" G_n and phase Φ_n can be defined for plasma density disturbances

$$G_n = \frac{2}{Rq_T} |n_0| = \frac{1 + \mu}{|\sqrt{1 + ik} + \sqrt{\mu^2 + i\beta k}|}, \quad \Phi_n = \arg(n_0) = \arg \frac{1}{\sqrt{1 + ik} + \sqrt{\mu^2 + i\beta k}}. \quad (11)$$

Characteristic length Λ_T of heated region along magnetic field can be estimated as

$$\otimes_r = -|\tau_0| \left(\frac{d|\tau|}{d\zeta} \right)_{\zeta \rightarrow 0}^{-1} = \left(\frac{2}{1 + \sqrt{1 + k^2}} \right)^{1/2}. \quad (12)$$

Similarly, field aligned scale of plasma density irregularities Λ_n can be estimated from

$$\otimes_n = -|n_0| \left(\frac{d|n|}{d\zeta} \right)_{\zeta \rightarrow 0}^{-1}. \quad (13)$$

In the long scale limit $k \rightarrow 0$ one can obtain that

$$G_r = 1, \quad G_n = \frac{1}{1 + \mu}, \quad \Phi_r = \Phi_n = 0, \quad \otimes_r = 1 \quad (14)$$

as in case of absence of drifts. It is obvious that drift does not affect disturbances with long enough scales along drift velocity.

In the short scale limit $k \rightarrow \infty$ there is significant dependence of "efficiencies" and characteristic scales of irregularities along the magnetic field on wave vector k :

$$G_r = \frac{1}{\sqrt{k}}, \quad G_n = \frac{1}{(1 + \sqrt{\beta})\sqrt{k}}, \quad \Phi_r = \Phi_n = -\frac{\pi}{4}, \quad \otimes_r = \sqrt{\frac{2}{k}} \quad (15)$$

"Efficiencies" of excitation of irregularities for short scales decrease as $k^{-1/2}$ as well as their characteristic lengths along magnetic field. Phases of both electron temperature and density disturbances become constant in quadrate to heating source.

Dependencies of above parameters on wave number k are shown in Figure 1. Contours of constant temperature and plasma density are shown in Figure 2 for some values of wave number k .

3.3 Constant drift gradient

Transport equations along the magnetic field for disturbances of electron temperature T and number density N in the coordinate frame connected to moving with plasma sources can be written taking into account constant drift gradient $dV_0/dx\{z\} = \text{const}$ in the form [Gurevich, 1978]

$$\begin{aligned} \kappa_{\parallel} \frac{\partial^2 T}{\partial z^2} - z \frac{dV_0}{dz} \frac{dT}{dx} - \delta v T &= -Q_T \delta(z), \\ D_{\parallel} \frac{\partial^2 N}{\partial z^2} - z \frac{dV_0}{dz} \frac{dN}{dx} - \gamma N &= - \left(Q_N - \frac{D_{T\parallel}}{\kappa_{\parallel}} Q_T \right) \delta(z) - \frac{D_{T\parallel}}{\kappa_{\parallel}} \left(z \frac{dV_0}{dz} \frac{\partial T}{\partial x} + \delta v T \right). \end{aligned} \quad (16)$$

Used notation is the same as in (1). It is supposed that irregularities are induced by heating Q_T and density Q_N sources that move with plasma drift velocity V_0 .

Un-dimensional equations for variables (2) where L_V (3) is substituted by

$$L_V = \frac{L_T}{\delta v} \frac{dV_0}{dz} \quad (17)$$

for harmonic solutions of (16) $(\tau, n) \propto \exp(ik\xi)$ can be written as

$$\begin{aligned} \frac{d^2 \tau}{d\zeta^2} - (ik\zeta + 1)\tau &= -q_T \delta(\zeta), \\ \frac{d^2 n}{d\zeta^2} - (i\beta k\zeta + \mu^2)n &= -(q_N - Rq_T) \delta(\zeta) - R(ik\zeta + 1)\tau \end{aligned} \quad (18)$$

Un-dimensional sources (4) and denotations (6) are used here.

Substitution $\zeta \rightarrow -\zeta$ to (18) leads to complex conjugate system of equations that means that $(\tau, n)(-\zeta) = (\tau, n)^*(\zeta)$. Considering region of $\zeta > 0$ one can rewrite (18) as

$$\frac{d^2 \tau}{d\zeta^2} - (ik\zeta + 1)\tau = 0, \quad (19)$$

$$\frac{d^2 n}{d\zeta^2} - (i\beta k\zeta + \mu^2)n = -R(ik\zeta + 1)\tau$$

with boundary conditions

$$\tau|_{\zeta \rightarrow \infty} = 0, \quad 2\Re\left(\frac{d\tau}{d\zeta}\bigg|_{\zeta \rightarrow 0}\right) = -q_T, \quad \Im \tau|_{\zeta=0} = 0 \quad (20)$$

$$n|_{\zeta \rightarrow \infty} = 0, \quad 2\Re\left(\frac{dn}{d\zeta}\bigg|_{\zeta \rightarrow 0}\right) = -(q_N - Rq_T), \quad \Im n|_{\zeta=0} = 0$$

where \Re and \Im stand for real and imaginary parts respectively. Relations for imaginary parts of τ and n at $\zeta = 0$ are followed from continuity conditions at the source.

Solution of equation (19) for τ with appropriate boundary conditions (20) is

$$\tau = -\frac{q_T}{2\Delta_\tau} \text{Ai}^*[v(0)] \text{Ai}[v(\zeta)], \quad v(\zeta) = (ik)^{1/3} \left(\zeta + \frac{1}{ik} \right) \quad (21)$$

where $\text{Ai}(z)$ is Airy function that does not diverge at infinity, and

$$\Delta_\tau = \Re\{\text{Ai}[v(0)]\Re\{(ik)^{1/3} \text{Ai}[v(0)]\} + \Im\{\text{Ai}[v(0)]\}\Im\{(ik)^{1/3} \text{Ai}[v(0)]\}\} \quad (22)$$

Substituting (21) into density equation (19) and solving it with appropriated boundary conditions (20) one can obtain

$$\begin{aligned} n = & -\frac{q_N}{2\Delta_n} \text{Ai}^*[w(0)] \text{Ai}[w(\zeta)] + \frac{Rq_T}{2\Delta_n} \text{Ai}^*[w(0)] \text{Ai}[w(\zeta)] \\ & \cdot \left[1 + \frac{\pi}{\Delta_\tau} \Re\{(ik)^{2/3} J \text{Ai}^*[v(0)] \text{Bi}[v(0)]\} \right] \\ & + \frac{\pi Rq_T}{2\Delta_n} \left(i \frac{1}{\Delta_n} \{i\beta k\}^{1/3} \text{Ai}^*[w(0)] \text{Ai}[w(\zeta)] \right. \\ & \cdot \Im\left\{\left(\frac{ik}{\beta}\right)^{1/3} J \text{Ai}^*[v(0)] \text{Bi}[w(0)]\right\} - \left(\frac{ik}{\beta}\right)^{1/3} \text{Ai}^*[v(0)] \\ & \cdot \left\{ \text{Bi}[w(0)] \int_{\zeta}^{\infty} d\zeta' v(\zeta') \text{Ai}[v(\zeta')] \text{Ai}[w(\zeta')] \right. \\ & \left. \left. + \text{Ai}[w(\zeta)] \int_0^{\zeta} d\zeta' v(\zeta') \text{Ai}[v(\zeta')] \text{Bi}[w(\zeta')] \right\} \right), \end{aligned} \quad (23)$$

$$w = (i\beta k)^{1/3} \left(\zeta + \frac{\mu^2}{i\beta k} \right)$$

where $\text{Bi}(z)$ is another Airy function. Here in (23)

$$\Delta_n = \Re\{\text{Ai}[w(0)]\Re\{i\beta k\}^{1/3} \text{Ai}^*[w(0)]\} + \Im\{\text{Ai}[w(0)]\}\Im\{i\beta k\}^{1/3} \text{Ai}^*[w(0)]\} \quad (24)$$

and

$$J = \int_0^{\infty} d\zeta' v(\zeta') \text{Ai}[v(\zeta')] \text{Ai}[w(\zeta')] \quad (25)$$

Using asymptotic expansions of Airy functions for $|z| \rightarrow \infty$

$$\begin{aligned} \text{Ai}(z) &\rightarrow \frac{1}{2\sqrt{\pi} z^{1/4}} \exp\left\{-\frac{2}{3} z^{3/2}\right\} \\ \text{Bi}(z) &\rightarrow \frac{1}{2\sqrt{\pi} z^{1/4}} \left(\exp\left\{\frac{2}{3} z^{3/2}\right\} - i \exp\left\{-\frac{2}{3} z^{3/2}\right\} \right) \end{aligned} \quad (26)$$

one can show that in the long scale limit obtained expressions (21) and (23) for τ and n respectively can be reduced to well known expressions for the case without drifts:

$$\tau_{k \rightarrow 0} = \frac{1}{2} q_T e^{-\zeta}, \quad n_{k \rightarrow 0} = \frac{1}{2} q_N e^{-\mu\zeta} - \frac{1}{2} R q_T \frac{1}{1-\mu^2} (e^{-\zeta} - \mu e^{-\mu\zeta}) \quad (27)$$

Suppose now for simplicity that there is no source of plasma density $q_N = 0$. In this case phases of temperature and density disturbances coincide with the phase of heating source due to relation $(n, v)(-\zeta) = (n, v)^*(\zeta)$.

"Efficiencies" of heating G_τ and excitation of density irregularities G_n are defined by (10) and (11), while characteristic lengths along magnetic field of temperature L_τ and density L_n irregularities — by (12) and (13) respectively. It can be shown that $G_\tau = L_\tau$.

It is obvious that in the long scale limit $k \rightarrow 0$.

$$G_\tau = L_\tau = 1, \quad G_n = \frac{1}{1+\mu}. \quad (28)$$

In this case drift does not affect development of irregularities.

In the short scale limit $k \rightarrow \infty$

$$G_\tau = L_\tau = -\frac{2}{\sqrt{3}} \frac{\text{Ai}(0)}{\text{Ai}'(0)} k^{-1/3}, \quad G_n = -\frac{2}{\sqrt{3}} \frac{\text{Ai}(0)}{\text{Ai}'(0)} \frac{1+\beta^{1/3}}{1+\beta^{1/3}+\beta^{2/3}} \quad (29)$$

It is taken into account here that

$$\frac{1}{\text{Ai}(0)\text{Ai}'(0)} \int_0^{\infty} du \cdot u \text{Ai}(u) \text{Ai}'(\beta^{1/3}u) = -\frac{1+\beta^{1/3}}{1+\beta^{1/3}+\beta^{2/3}} \quad (30)$$

Let us also note that

$$\frac{2}{\sqrt{3}} \frac{\text{Ai}(0)}{\text{Ai}'(0)} = \frac{2\Gamma(1/3)}{3^{5/6}\Gamma(2/3)} \approx -1.58393 \quad (31)$$

where $\Gamma(x)$ is Gamma function. "Efficiencies" of excitation of irregularities for short scales decrease as $k^{-1/3}$ as well as their characteristics lengths along magnetic field.

Dependencies of "efficiencies" G_τ and G_n on wave number k are shown in Figure 3.

3.4 Conclusions

Let us estimate relevance of above considerations to the Earth's ionosphere. Typical values of ionospheric parameters in the F-layer are $(\delta v)^{-1} = 10-30$ s, $L_T = 30-100$ km, and drift velocity at mid latitudes can vary in the range of $V_0 = 10-200$ m/s. As it can be seen from Figures 1, 3 drifts can essentially affect development of irregularities with scales of order or less than L_V for fixed heating source and L_V' for a source moving with plasma drift.

In case of ionospheric modification by powerful radio waves fixed source is realized for narrow beams, for radio waves focused by natural horizontal inhomogeneities in lower layers of ionosphere, and for self-focusing of powerful radiation. Typical values of characteristic scales of irregularities when drift influence becomes significant are $L_{\nu} = 100-6000$ m for parameters given above. These scales lay in typical interval both for natural focusing and for self-focusing of radio waves.

Drift velocity gradient becomes essential when assumption of homogeneous electric field of powerful radio wave is applicable for description of excitation of irregularities as in case of resonance instability or thermal parametric instability. Gradient of drift velocity can exist in the ionosphere even just due to the dependence of magnetic field strength on the altitude. Drift velocity in crossed electric E and magnetic B fields is described by relation $V_0 = cE/B$. Supposing that magnetic field lines are equi-potential one can found that $E \propto S^{-1/2}$ where S is area of magnetic flux tube, $S \propto 1/B$. Hence, $V_0 \propto B^{-1/2}$. Considering Earth's magnetic field as dipole field with $B \propto R^{-3}$ one can obtain for drift velocity $V_0 \propto R^{3/2}$ where R is distance from the center of the Earth, and $dV_0/dz = 3 V_0/2R$. Gradient of drift velocity in the F-region of the ionosphere can be estimated as $0.2-4.5$ sm/s /km for above parameters. Characteristic scales for drift gradient $L_{\nu'} = 0.6-150$ m are comparable to those usually considered in theories of small scale stratification of ionospheric plasma by powerful radio waves. Drift gradient can also essentially curve irregularities from the straight line along magnetic field and determine aspect sensitivity of radio wave scattering off.

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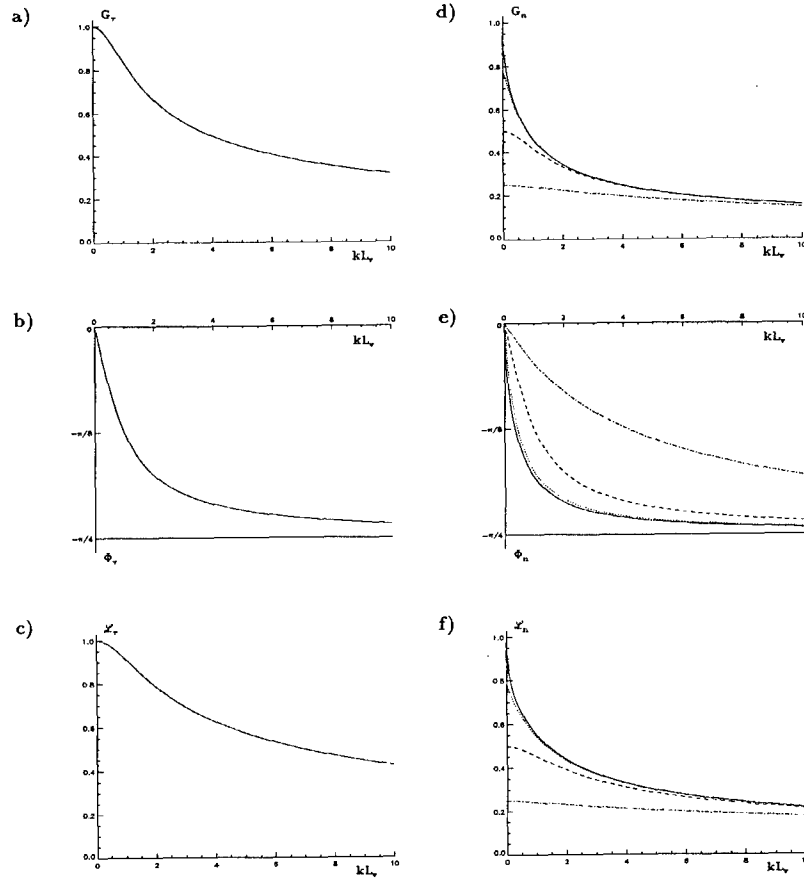


Figure 3.1: "Efficiencies" of heating G_{τ} (a) and generation of irregularities G_n (d), phases of heating Φ_{τ} (b) and irregularities Φ_n (e), and characteristic lengths along magnetic field of electron temperature Λ_{τ} (c) and density irregularities Λ_n (f) versus transverse scales of irregularities. In panels (d) , (e), and (f) curves are shown for $\beta=1$ and $\mu=0$ (solid lines), 0.3 (dots), 1 (dashes), 3 (dash-dots).

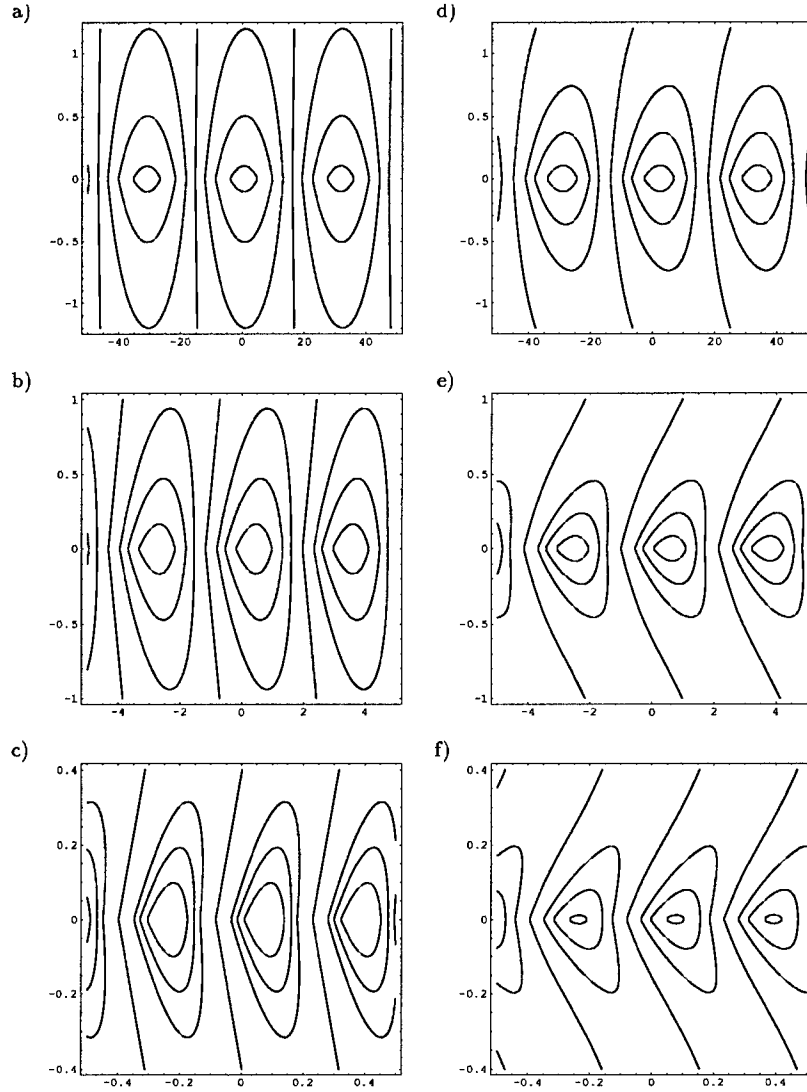


Figure 3.2: Contours of constant electron temperature / (a), (b), (c)/ and plasma density / (d), (e), (f)/ for $\beta=1$ and $\mu=0$ for different values of wave number k . Along horizontal axis: $\xi = x/L_V$, along vertical axis: $\zeta = z/L_T$. Panels (a) and (d) correspond to $kL_V = 0.1$, absolute contour values for panel (a) are 0.9, 0.6, 0.3, 0, for panel (d) — 0.7, 0.5, 0.3, 0. Panels (b) and (e) correspond to $kL_V = 1$, absolute contour values for panel (b) are 0.7, 0.5, 0.3, 0, for panel (e) — 0.4, 0.3, 0.2, 0. Panels (c) and (f) correspond to $kL_V = 10$, absolute contour values for panel (c) are 0.25, 0.2, 0.15, 0, for panel (f) — 0.15, 0.11, 0.07, 0.

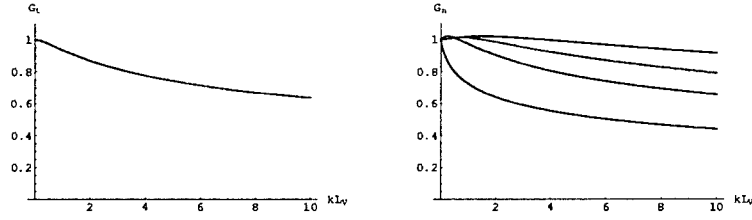


Figure 3.3: "Efficiencies" of heating $G_t = \Lambda_t$ (a) and generation of irregularities G_n (b) versus transverse scales of irregularities. Curves in panel (b) are shown for $\beta = 1$ and $\mu^2 = 0, 0.3, 1, 3$ (lower curve corresponds to $\mu=0$).

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Interim Report

The work is driven by the intersection of two lines of activity, one experimental, the other theoretical, that converge on an importantly new view of the physics of how high power HF radio waves interact with collisional plasmas, such as the earth's ionosphere.

Its goal is to seize this opportunity to amalgamate the complementary US and Russian recent findings, to strengthen the foundation of plasma physics underlying future work in the growing field of ionospheric HF heating. The following main problems are considered in the present Interim Report:

1. Observations of fast electrons in a wide region around O-wave reflection point and comparison of observational data with multiple acceleration theory.
2. Energization of electrons by Langmuir turbulence in the plasma resonance region and its relation to the observations of excited optical emission.
3. Drift effects on the scale of field-aligned irregularities.

Publications

1. A. V. Gurevich H. C. Carlson, G. M. Milikh, K. P. Zybin, F. T. Djuth and K. Groves Suprathermal electrons generated by the interaction of powerful radio wave with the ionosphere, GRL 2000 (accepted for publication).
2. A.V. Gurevich, H.C. Carlson, G.M. Milikh and K.P. Zybin Optic emission from modified by powerful radio waves ionosphere (in preparation).

1. Experimental Studies of Energetic Electrons in Ionosphere Using ISR Technique and Comparison of Observational Data with Multiple Acceleration Theory

Abstract

Two sets of observations of suprathermal electrons, produced by the interaction of powerful radio wave with the ionosphere, as seen using the incoherent scatter radar (ISR) technique are presented. The observational data are compared with the theory of multiple acceleration of electrons in the strongly excited resonance region near the reflection point of the powerful radio wave. The structure of the wide perturbed region filled with energetic (10–20) eV electrons is determined. The size of this region along the Earth magnetic field is shown to be 100 km. The full power going to the accelerated particles is determined: it is 6–8 kW or (4–6)% of the entire HF radiated power. The power carried by the suprathermal electrons flux escaping into magnetosphere is of the order 1 kW.

1.1 Introduction

Observations of fast suprathermal electrons, due to the interaction of powerful radio waves with the ionosphere, using incoherent scatter radar (ISR) technique were first proposed and realized by Carlson, Wickwar and Mantas in 1972 (see Carlson et al. [1982]). The powerful O-mode wave in the resonance region near the reflection point effectively excites plasma waves – natural oscillations of the ionosphere plasma. Plasma waves due to the nonlinear cavitation process accelerate suprathermal electrons to energies in excess of $\varepsilon \sim 10\div 20$ eV, i.e. two orders of magnitude higher than the thermal electron energy $T_e \approx 0.1$ eV. The fast electrons propagate in the ionosphere to large distances, 10–100 km from the acceleration region. Those electrons excite nonthermal plasma waves, which were detected by ISR [Carlson et al., 1982].

We discuss here the results of two sets of experiments. The experimental data are here compared in detail with theory based on the mechanism of multiple electron acceleration.

Carlson et al. [1982] gave a qualitative explanation of their observations. They pointed out the significant role played by multiple scattering of electrons in the neutral atmosphere. The quantitative theory of multiple acceleration was proposed by Gurevich et al. [1985]. The comparison of the theory and observations can help reach a still higher level in our understanding of a number of space plasma phenomena in which accelerated electrons play an important role.

1.2 Observations

The most sensitive and direct ground based means of detecting 10–20 eV suprathermal electrons is by incoherent scatter plasma line observations of the Langmuir waves produced in the ionosphere by these electrons [Perkins and Salpeter, 1965]. This technique has been used to study photoelectron fluxes in the ionosphere [Yngvesson and Perkins, 1968]. Carlson et al. [1982] applied the ISR technique to measurements of suprathermal electrons produced by the interaction of powerful radio wave with the ionosphere at Arecibo. We present here two types of experimental data. The primary and secondary data were gathered on the nights 20 May 1972 and 13 July 1992 respectively. The primary data were published by Carlson et al. [1982], here we will revisit those results and supplement them with the complimentary set of secondary 1992 data [Djuth et al., 1996].

For the 1972 Arecibo data here, a HF dipole feed hung over the 1000-foot diameter ISR radar dishes, for which the 138 kW transmitter at 7.63 MHz delivered about 30 W/m² to the HF reflection height. For the 1992 Arecibo data here, a separate 48 array of log periodic antennas 17 km away from the ISR radar dish, for which the 300 kW transmitters at 5.1 MHz, into 50% efficiency transmission lines, delivered about 40 W/m² to the HF reflection height.

For 20 May 1972, strong plasma waves were excited near the reflection point at the altitude of $z = 285$ km. Night time plasma line intensities were observed to be enhanced by a factor of 10–100, over a range of altitudes which extended to below 250 km. They relaxed to their normal level within the electron transport time after the transmitter was turned off.

The 430 MHz incoherent scatter radar was used to diagnose the ionosphere, which allowed measurement of altitude profiles of the background plasma temperature T_e and electron concentration N . The plasma line echo from the plasma waves excited by the suprathermal electrons provided the altitude of a set of plasma frequencies f_p between 5 and 7.5 MHz, with an altitude resolution of 1.5 km.

For a given diagnostic radar wavelength λ_r , the main scattered signal comes from ionosphere plasma waves, whose wave vector is directed toward the radar, and the phase velocity is equal to the velocity of suprathermal electrons v

$$v_{ph} = \frac{1}{2} \lambda_r f_p \quad (1)$$

Here $f_p = (e^2 N / 4\pi m)^{1/2}$ is the local plasma frequency. For the Arecibo radar $\lambda_r = 70$ cm and from eq. (1) one can obtain that the energy of electrons is

$$\varepsilon = \frac{1}{2} m v_{ph}^2 = 0.35 f_p^2 V \quad (2)$$

The intensities of plasma line echoes were shown in Fig.1 of Carlson et al. [1982]. A significant enhancement of radar echo intensities is evident at large distance below the reflection point of the radio wave at 285 km. The upgoing one-dimensional fluxes of suprathermal electrons at the heights 266 km and 256 km determined from these data are shown in Fig.3 of Carlson et al. [1982].

The additional data were gathered on the night 13 July 1992. The reflection altitude was 295 km, with strong enhancements of plasma line intensity observed in the vicinity of the reflection level. An example of plasma line echoes collected at the heights 337–367 km, where the ISR beam intersected the flux of the accelerated electrons, is shown in Fig.1.1. The echoes were obtained both below and above the maximum of F-layer. Electron density $N_{max} \approx 5.5 \cdot 10^5 \text{ cm}^{-3}$ was reached at F-maximum heights $z = 350$ km. The observed plasma line echoes correspond to upgoing electron flux. Note particularly here a significant enhancement of plasma line echoes even at altitudes more than 70 km above the reflection layer.

1.3 Brief Outline of the Theory

In ionosphere modification experiments, electrons gain energy in a strongly disturbed Langmuir resonance layer near the reflection point of the O-mode wave. The essential point of multiple acceleration theory is a large fraction of fast electrons, after leaving the acceleration layer will return back due to collisions with neutral particles, and thus can gain additional energy. The process can be repeated many times. Thereby, a wide region around the acceleration layer, which is elongated along the geomagnetic field, can be filled with strongly heated magnetized suprathermal electrons. The elongation along the

Earth's magnetic field can reach an order of a hundred kilometers because it is determined by two factors: the large mean free path of the suprathermal electrons and a small part of electron energy is lost in one collision. The extent of upward elongation will be considerably greater than in the downward direction, because the mean free path above the reflection height is considerably greater than below it. Consequently, a significant flux of suprathermal electrons moves to the magnetosphere, and can even reach the magnetic conjugate point.

The acceleration process due to the multiple crossing of acceleration layer is averaged and in final form, depends on two scalar factors only: the full power density P absorbed by fast electrons in the acceleration layer, and a characteristic parameter describing the effectiveness of the acceleration inside the layer T_{ef} . They are related to the details of the accelerating process, effective number of cavitons, their width, and so on.

The distribution function of suprathermal electrons in the theory of multiple electron acceleration can be presented in the simple form [Gurevich et al., 1985]

$$f_0(\varepsilon, z) = CK_0\left(\frac{\varepsilon}{T_{ef}}\right) \exp\left\{-\left|\int_0^z \frac{dz}{L_\varepsilon^\pm \cos \alpha}\right|\right\} \quad (3)$$

Where K_0 is a modified Bessel function, ε is the electron energy and T_{ef} is the effective temperature of suprathermal electrons, while α is the angle between the vertical and geomagnetic field. The normalization constant C is directly proportional to the power density P of the HF wave absorbed by the suprathermal electrons:

$$C = \frac{m^2}{4\pi^3 T_{ef}^3 \left(\frac{\delta}{3}\right)^{1/2}} P \quad (4)$$

Here δ is the average fraction of electron energy lost in a single collision with neutral molecules.

As given in eq. (3) the acceleration layer is assumed to be located at $z = 0$, while the L_ε^\pm factor is the characteristic relaxation length of suprathermal electrons in the upward (+) and downward directions (-)

$$L_\varepsilon^\pm(z) = \left[N_m^\pm \sigma_{tr}(\varepsilon) (3\delta)^{1/2} \right]^{-1} \quad (5)$$

Where $N_m^\pm = N_m^\pm(z)$ is the neutral density above (+) and below (-) the layer, $\sigma_{tr}(\varepsilon)$ is the total transport cross-section of electron-neutral collisions, and $\delta = \sigma_{in} / \sigma_{tr}$, where σ_{in} is the total cross-section of inelastic collisions, which includes ionization by electron impact. In σ_{tr} and σ_{in} the collisions with all neutral components are considered.

1.4 Comparison of the Observations with Theory

In the experiment on 20 May 1972, upgoing (in the radar look direction) fluxes of suprathermal electrons in the energy range $10 \div 17$ eV at the heights $z_1 = 256$ km and $z_2 = 266$ km were determined. The reflection point was at $z_0 = 285$ km. Comparison with the theory requires determination of the electron flux along radar direction e_r :

$$J_r = (\mathbf{v}, \mathbf{e}_r) f(\mathbf{v})$$

We introduce net upgoing and downgoing fluxes by integrating the flux J_r over corresponding angles in velocity space. Taking the angle between the geomagnetic field and vertical radar ray in Arecibo as $\alpha = 40^\circ$ we obtain the upgoing J_+ and downgoing J_- fluxes in energy interval $d\varepsilon$:

$$J_{\pm} = \frac{2\pi}{m^2} \mathcal{E}_0(\varepsilon) \left(1 \pm \frac{2}{3} (3\delta)^{1/2} \cos \alpha \right) \exp \left(- \left| \int_0^z \frac{dz}{L_{\varepsilon}^{\pm}(z) \cos \alpha} \right| \right) \quad (6)$$

The results of calculations are presented in Fig.1.2 for the given heights and different T_{ef} . It is apparent from the figure, that in the energy range $10 \div 20$ eV reveal a rather flat spectrum. Besides, there is no strong difference between the fluxes J at different temperatures T_{ef} .

On the other hand, as follows from calculations the fluxes are effectively diminishing with distance from the acceleration layer. This is because electron energy is lost in inelastic collisions with the neutral. Equation (7) implies that the fluxes of upgoing electrons are smaller than the downgoing fluxes below the acceleration region. Their relation depends upon the angle α between radar ray direction and the geomagnetic field – for Arecibo at $\alpha = 40^\circ$ $J_+/J_- \approx 1.5$. Comparison between the theory and observations presented in Fig.1.3 shows a reasonable agreement between those two. In fact, the behavior of the spectrum of the electron flux at different heights is consistent with the theory for any of T_{ef} applied. Taking into account the absolute values of upgoing flux $J_+ = (4 \div 8) \times 10^5$ el/cm²s eV at characteristic energies $10 \div 15$ eV obtained by Carlson et al. [1982], one can find the absorbed power W_S of the HF wave converted into acceleration of the suprathermal electrons

$$\begin{aligned} \text{for } T_{ef} = 5 \text{ eV} \quad W_S &\approx 5.6 \text{ kW} \\ \text{for } T_{ef} = 7.5 \text{ eV} \quad W_S &\approx 6.2 \text{ kW} \\ \text{for } T_{ef} = 10 \text{ eV} \quad W_S &\approx 8.5 \text{ kW} \end{aligned}$$

In these calculations the Arecibo heater beam was taken roughly 10^3 km² for a full absorbed power $W_S = PS$. The factor P was determined from calculations of the electron flux made at different distances from the source, and from the results presented in Fig.1.3. It is apparent that the dependence of W_S on T_{ef} is not very strong. Note that eq. (7) allows us to estimate the upgoing electron flux at $z > 500$ km which escapes to the magnetosphere due to the absence of scattering collisions.

Figure 1.4 shows the height distribution of plasma line intensity measured during the second experiment, which is compared with the height dependence of the distribution function of suprathermal electrons $f_0(\varepsilon(z), z)$. Here energy $\varepsilon(z)$ is determined through the measured electron density distribution (Fig.1.1) by using eq. (2). One can see a sufficient agreement between the theory and observations. Note, that the normalized height dependence of the distribution function practically does not depend on the parameter T_{ef} (see Fig.1.4).

1.5 Discussion and Conclusions

It was shown that the ISR plasma line measurements are in reasonable agreement with the theory of multiple electron acceleration. From the observational data and their comparison with the theory it follows that: In experiments where powerful radio waves interact with the ionosphere a large number of suprathermal electrons are generated in the energy range up to 20 eV.

The suprathermal electrons are observed over a wide altitude region of the order of a few tens km both below and above the acceleration layer, centered near the reflection altitude of the powerful radio wave. While not observed, they must extend well above this altitude, into the plasmasphere/magnetosphere, and even to some degree into the magnetically conjugate region.

In power, $W_S \approx 5.6 \div 8.5$ kW ($4 \div 6$)% of the full transmitted heater power goes into the acceleration of suprathermal electrons. This energy is dissipated due to the

generation of optical emissions, ionization and heating of ionized and neutral components of the ionosphere plasma. The characteristic dissipation length for suprathermal electrons depends strongly on the altitude. For $z = 250$ km it is about 10 km, while for $z \approx 300$ km it is about 30 km.

The flux of suprathermal electrons into the plasmasphere/magnetosphere depends strongly on the height z_0 of the acceleration layer. For the studied cases $z_0 = 285$ km, and $z_0 = 295$ km, and $W_S = 8.5$ kW, $T_{ef} = 10$ eV the flux into magnetosphere in the energy range $10 \div 20$ eV is about 0.9 kW. This estimate is consistent with earlier calculations of photoelectron energy and flux escape into and loss to the plasmasphere/magnetosphere [Mantas et al., 1978].

We conclude that plasma line observations combined with the theory enables us to obtain significant information about the acceleration of the suprathermal electrons in ionosphere modification experiments. For the first time the structure and size of the perturbed region filled with suprathermal electrons along with the full power going into the accelerated electrons is determined by a theoretical calculation consistent with an observational data set.

Note that we had the possibility to use here only a small part of ISR data on ionosphere existing from experiment 1992. The further elaboration of this data and their comparison with the theory is of a significant interest and will be definitely quite fruitful.

We suggest also these experiments be repeated to obtain more information about the main features of acceleration, and its dependence upon the reflection height and ionosphere conditions. Note that such observations have as yet been done only at Arecibo, though we expect that the latitude dependence of the effect could be very significant.

At the same time, the agreement that found between the theory and observations make it possible to plan more detailed future experiments. These will thus help to reach a much better understanding of the physical mechanisms of electron acceleration and nonlinear processes in ionosphere plasmas.

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Figure Captions

Figure 1.1 Spectra of the downshifted plasma line measured at Arecibo 13 July 1992 at 8:30 LT.

Figure 1.2. Upgoing flux of suprathermal electrons versus their energy computed for the full absorbed power $W_S = 10$ kW and for different values of T_{ef}

Figure 1.3. Upgoing flux of suprathermal electrons versus their energy. Solid curves 1, 2 and 3 correspond to $T_{ef} = 5, 7.5$ and 10 eV respectively, while the points with bars correspond to the observations made on 052072.

Figure 1.4. Plasma line intensities versus altitude. Solid curves 1, 2 and 3 correspond to $T_{ef} = 5, 7.5$ and 10 eV respectively, while the points with bars correspond to the observations made on 071392.

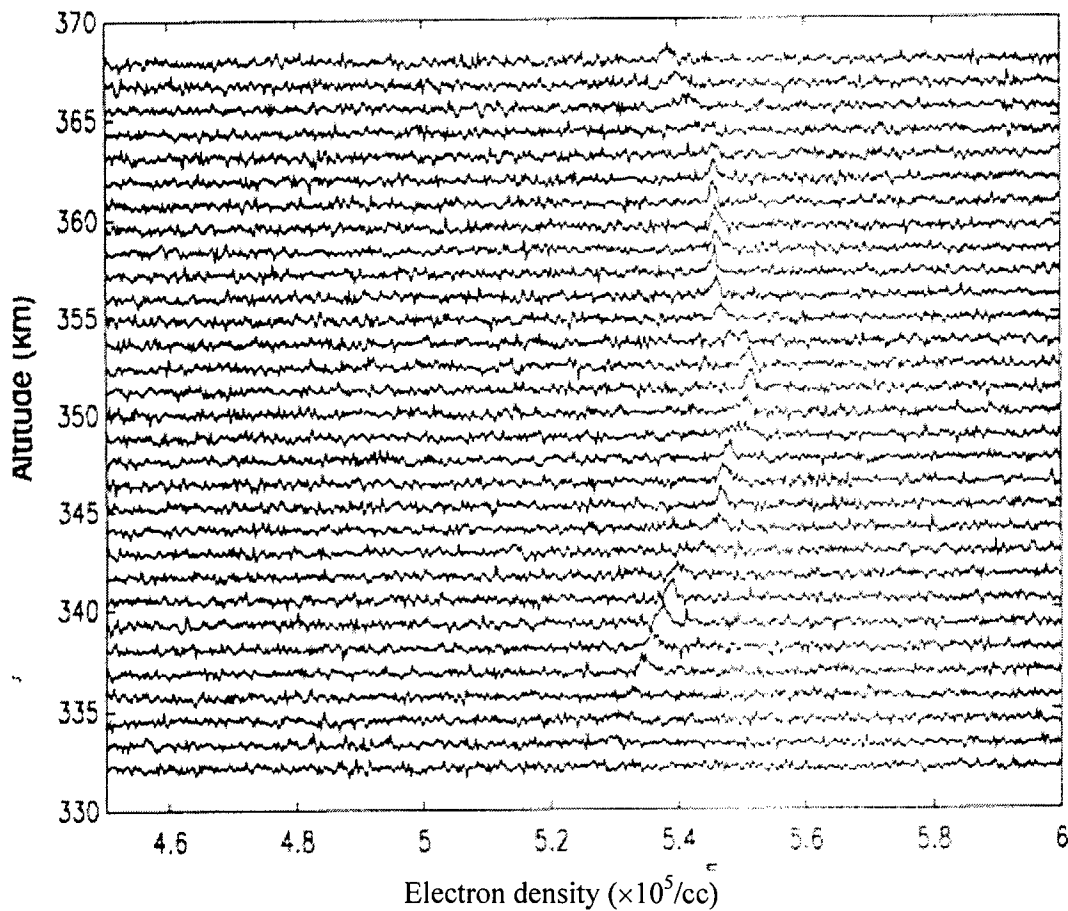


Figure 1.1

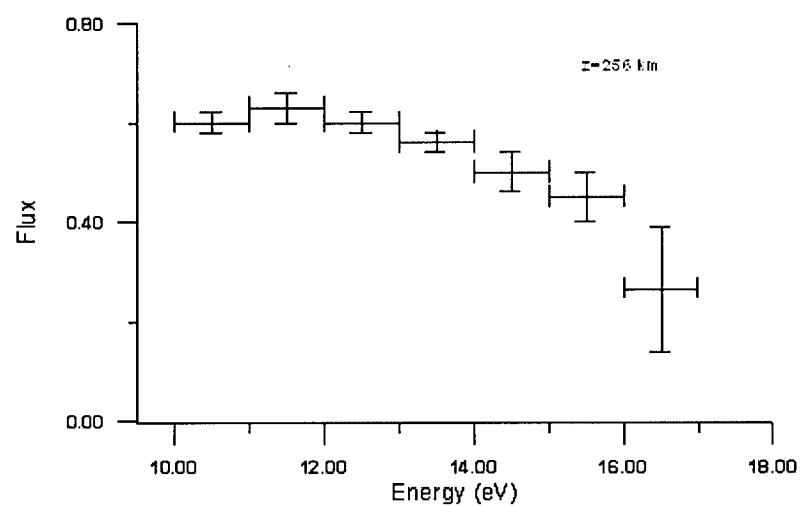
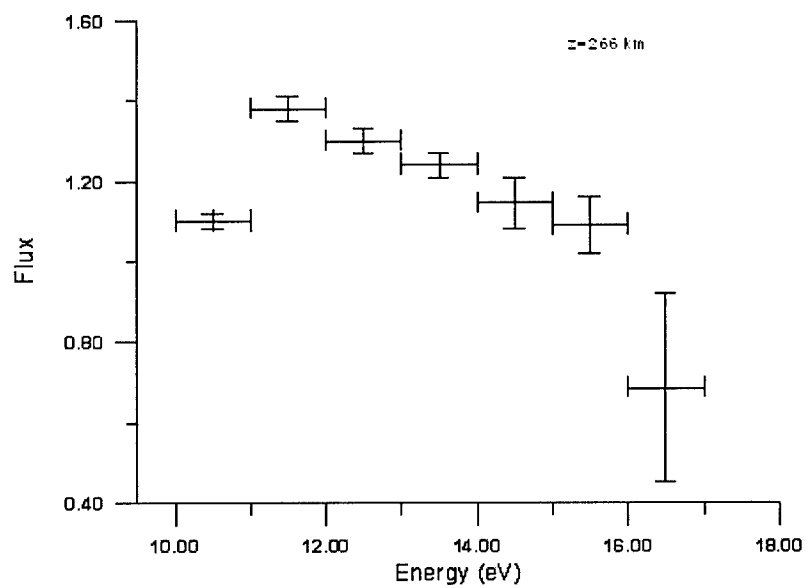


Fig.1.2

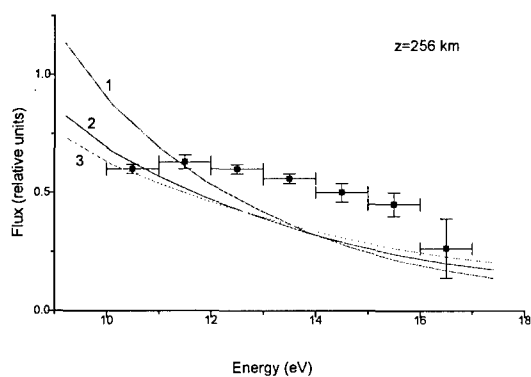
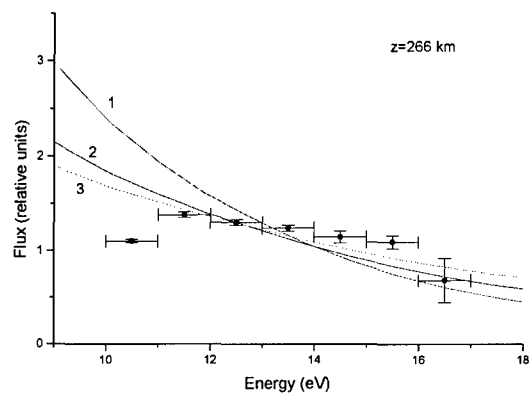


Figure 1.3

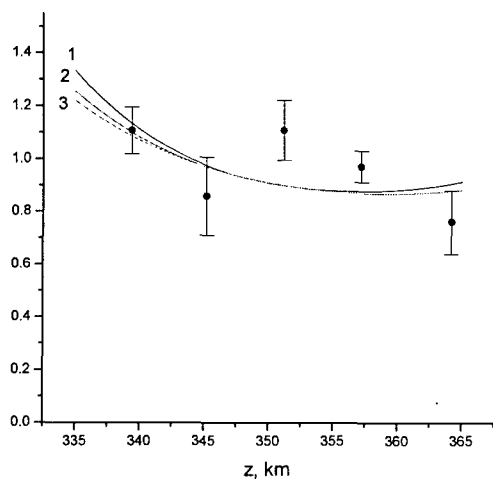


Figure 1.4

2. Optic emission from modified by powerful radio waves ionosphere

Introduction

Enhancement of optic emission was established in the first ionosphere modification experiments (Sipler and Biondi, 1971; Adeishvili et al., 1976). It was studied afterwards in multiple observations (Haslett and Megill, 1974; Carlson, 1974; Bernhardt et al., 1989). During the last years new interesting results have been obtained at HAARP facility (Peterson, 1999; Carlson, 1999), in sporadic E-layer at Arecibo and at low duty circle experiments at SURA (Nazyrov et al., 1999). The optic emission measurements indicate directly the existence of suprathermal electrons accelerated up to the energies 10 eV and more in modified ionosphere.

Other type of observations, demonstrating the accelerated up to the energies 20 eV electrons was performed by Carlson et al., 1982, Juth et al., 1996, using incoherent scattering radar at Arecibo. The suprathermal electrons were seen in these experiments in a wide region of the order 100-km around the reflection point of the heater wave. Some indications of the significant effect of suprathermal electrons on SEE emission exist as well (Frolov, 1998).

Theoretical considerations from the very beginning connected the enhancement of optic emission with the acceleration of electrons in the Langmuir plasma turbulence layer near the reflection point of powerful HF O-wave. Though, it was indicated as well, that the mostly intensive red-line emission (630 nm) partly could be determined by the direct effective heating of the main bulk of electrons (Carlson, Mantas, 1996), or it's significant part, captured inside nonlinear density depletions – "striations" generated by the upper-hybrid turbulence (Gurevich and Milikh, 1998). The theory of Langmuir plasma turbulence generated near the reflection region of pump O-wave connects the electron acceleration process with the modulation instability and with the creation of nonlinear Langmuir cavitons – density depletions filled with Lagmuir plasma oscillations (Wong and Stenzel, 1975). The fast electrons, passing trough the turbulent region filled with cavitons obtain the energy from the trapped caviton plasma oscillations in a diffusive way (Eastbrook et al., 1975; Morates and Lee, 1977; Andreev et al., 1980; Wang, Goldman and Newman, 1997). The main feature of this mechanism is that the energy is gained only by fast electrons whose velocity is high enough $v > a/\omega$, where a is a scale of a caviton, ω is the heater wave plasma frequency. Low energy electrons oscillate in the caviton adiabatically and do not get any additional energy. It should be emphasized, that the collapse of cavitons does not effect strongly the acceleration process of fast electrons (Wang et al., 1997). In the same time collapse can input a significant part of wave energy to the thermal component (Du Bois et al. 1993, Hanssen et al. 1997). To calculate the flux of energy, which goes to fast electrons the detailed, kinetic theory is needed, which takes into account electron collisions.

For example, the number and energy distribution of fast electrons in ionosphere conditions is affected by the so-called "multiple acceleration" (Gurevich et al., 1985, Vaskov et al., 1983). The multiple acceleration process is determined by the fact, that the Langmuir acceleration take place in a strongly excited layer, embedded in a weakly disturbed ionosphere plasma. After the fast electrons leave the acceleration layer it collides with a neutral molecules in nondistrubed plasma. Due to collisions some electrons could return to the acceleration layer and get there the energy once more. The process could be repeated many times, thus creating a "multiple acceleration". Directly this process determines the high-energy tail of electron distribution function and a wide spread of fast electrons around the acceleration region.

Though, understanding of the main mechanisms of the acceleration process in Langmuir turbulence seems significantly clear enough, both the theory and observations in ionosphere modification experiments up to now were not developed to the state when a detailed quantitative comparison between them could take place. That means, that some fundamental problems of modification experiments remain non-answered. One of them is the intensity and space distribution of the observed optic emission.

In the present work we intend to develop the existing theory and apply it to obtain the detailed comparison with the main part of the existing experimental data. The paper is constructed as follows. In section 2 heating of electrons by both in upper hybrid and in Langmuir turbulence is considered. It is demonstrated that the calculations of electron acceleration in Langmuir resonance layer without electron collisions lead to controversial result. In section 3 the theory of multiple acceleration is generalized. The energy losses of electrons on the excitation of molecular vibration level are taken into account. In section 4 the model calculations describing the general peculiarities of the suprathermal electron space energy distribution are performed. The section 5 describes the detailed comparison of the theory with the main observational data of the artificial optic emission in F-layer. The strong enhancement of the optic emission in sporadic E-layer modification is discussed in sections 6. In conclusion we formulate the nowadays state of the problem. The strongly inhomogeneous character of the disturbed region is stressed.

1.2 Heating of electrons in the resonance layer

In the vicinity of reflection point of powerful O-wave in ionosphere both heating and acceleration of plasma electrons takes place.

As is well known for vertically propagating pump O-wave the resonance region is situated at the heights Z between upper hybrid Z_{uh} and Langmuir Z_L resonances

$$Z_{uh} \leq Z \leq Z_L \quad (1)$$

In this region the frequency of the pump wave ω coincide with the frequency of the natural oscillations in ionosphere plasma

$$\omega = \omega_L = \sqrt{\frac{4\pi e^2 N(Z_L)}{m}} \quad (2)$$

Here $N(Z_L)$ – electron density at the Langmuir resonance, where the Langmuir turbulence take place. And Z_{uh} is defined by resonance condition for upper hybrid resonance

$$\omega = \omega_{ph} = \sqrt{\omega_L^2 + \omega_c^2} \quad (3)$$

In the vicinity of Z_{uh} the upper hybrid (UH) turbulence is effectively developing. The width of resonance region ΔZ depends on the pump wave frequency and plasma density gradient L :

$$\Delta Z = Z_L - Z_{uh} = L \frac{\omega_c^2}{2\omega^2} \quad L = \frac{N}{\left| \frac{dN}{dZ} \right|} \quad (4)$$

ω_c is electron gyrofrequency. For the ionosphere F-layer conditions usually $\Delta Z \sim 2 \div 5$ km.

Conventional collision absorption of radio waves in F-region is very weak – less than 1 dB [Ginzburg, 1967]. If plasma turbulence in the resonance region is excited the absorption become abnormally strong – up to 10÷20 dB [Gurevich, 1978].

Note, that the Langmuir turbulent layer (LT) is placed in the upper part of resonance region close to the reflection point (2), while close to the upper hybrid resonance (3) lays upper hybrid turbulent layer (LH). Acceleration of electrons take place

in Langmuir resonance layer (we do not speak here about singled out case of multiple gyroresonance frequency $\omega = n\omega_c$). The UH turbulence lead to formation of striations. An effective anomalous absorption of the pump wave take place in this region, which lies below the Langmuir turbulence layer. It means that the UH anomalous absorption (absorption on striations) can affect significantly the heating and acceleration processes in the Langmuir turbulence layer while diminishing the pump-wave power.

Heating of electrons in UH turbulence region

Anomalous absorption is determined by the excitation of UH waves due to direct transformation of the pump wave on the density gradient in striations. UH waves trapped in striations heat electrons and thus determine the depth form and number density of striations. Nonlinear theory of the electron heating and anomalous absorption on striations was developed by Gurevich et al. 1996. The results of the theory are in a rather good agreement with observations.

Following to the theory of UH anomalous absorption we determined the simple relation for the losses of pump wave power P in the UH turbulence region on its way to Langmuir turbulent layer.

$$P(Z_L) \approx P_0 10^{-K/10} \quad K = \gamma \left(\frac{L}{100 \text{ km}} \right) Q$$

$$\gamma \approx 5.5 \div 8, \quad 0 \leq Q \leq 1$$

Here K is an attenuation coefficient in dB. Coefficient Q is proportional to the relative depth of striations. It depends on pump wave power, frequency, plasma parameters. The state of the striation development is also significant. For the well developed striations $Q \sim 1$.

1.3 Heating of electrons in Langmuir turbulent layer

The excitation of Langmuir turbulence, its steady state and caviton formation in the field of powerful radio wave in conditions close to ionosphere was studied numerically. We will discuss the result of the calculations of Hanssen et al and Wang et al in order to determine the thermal heating of electron by Langmuir turbulence. The parameters used in calculations are presented in Table 1.

	Hanssen et al (1992)	Wang et al (1997)
Pump wave frequency (MHz)	f= 4.01	f=4.9
Pump wave field E_{rms} (V/m)	1.26	1.6
Plasma density n_0 (cm ⁻³)	2×10^5	2×10^5
Electron temperature (eV)	0.17	0.1
Pump wave energy $W_{rms} = E_{rms}^2 / 4\pi m_0 T_0$	2.7×10^{-3}	5×10^{-3}

Hanssen et al. (1992) considered one-dimensional Zakharov system of equations describing nonlinear development of plasma density and electric field fluctuations under the action of homogeneously oscillating electric field of the pump wave. Electron and ion collisions and Landau damping were taken phenomenologically, using terms from linear theory. The established average energy of longitudinal plasma waves, concentrated mostly in cavitons grew up 70 times in comparison with the pump wave energy. It gives $W_{rms}^p \approx 0.19$.

The power density dissipating by plasma oscillations to the heating of electrons is

$$\frac{dW}{ds} = \frac{e^2 E_{rms}^2 v_e n_0}{m \omega^2} \quad (5)$$

Where s is coordinate along the Earth magnetic field, v_e – electron collision frequency, ω – plasma wave frequency and $E_{rms}^2 = 4\pi n_0 T_e W_{rms}$. Following Hanssen et al. let us suppose that the whole width of the turbulent layer is of the order of pump wave length near reflecting point $L \sim \lambda \approx 200$ m. We obtain that the full pump wave power dissipating in Langmuir cavitons to the heating of plasma electrons:

$$W_T \approx W_{rms} v_e T_e \lambda_m \approx 40 \frac{\mu kW}{m^2} \quad (6)$$

That value is quite significant. Really, if we neglect anomalous absorption in striations below turbulent layer, then to create the electric field $E_{rms} = 1.26$ V/m (see Table 1) at the main Eiry maximum near reflection point, the pump wave power $60(\sin\alpha)^{4/3} (100km/L)^{1/3} \mu KW/m^2$ is needed. Here α is the inclination angle of Earth magnetic field. We see that at $\alpha \sim 1$ about 5 dB of pump wave is absorbed due to the thermal heating of electrons in the cavitons. Note that the calculated pump wave absorption is effectively growing with latitude ($\alpha^{4/3}$). So, more than 60% of absorbed pump-wave energy in caviton turbulence is going directly to the thermal component of electrons.

Wang et al. 1997 elaborated one-dimensional kinetic model. Not only field equation but also Vlasov kinetic equation for noncolliding electrons was integrated simultaneously in the model. The Langmuir waves instability was excited and established. Energy level of average plasma oscillations grew up about 30 times in comparison with pump wave field reaching $W_{rms} = 0.14$. We see that in spite of quite different approach, the numerical value of saturated energy of plasma oscillations is close enough in both calculations.

The power density dissipating to the thermal component of electrons is again determined by expression (6). The width of turbulent layer was supposed to be $\lambda_m \approx 180$ m. From (6) we again obtain the heating power $W_T = 40 \mu kW/m^2$, what means that about 2.5 dB of pump wave energy goes to the heating in cavitons.

We see that when Langmuir turbulence is effectively excited the pump wave energy losses to the heating of electrons is quite significant, compatible (though remaining less) to the anomalous absorption in UH layer.

Significant advantage of Wang et al. 1997 kinetic model is the possibility to see directly not only saturation of plasma oscillations, but acceleration of suprathermal electrons as well. The result of calculations shows that a tail of initial Maxwellian distribution function is growing in time. An established suprathermal stationary tail is

$$f_s = 8.3 \times 10^{-6} \left(\frac{32V_{Te}}{v} \right)^{1.9} \frac{n_0}{V_{Te}} \quad 4 \leq \frac{v}{V_{Te}} \leq 25 \quad (7)$$

shown to be approximated by power – law:

Here $V_{Te} = (T_e/m)^{1/2}$ – thermal velocity. The one dimensional maxwellian part of distribution function has standard view:

$$f_m = \frac{n_0}{\sqrt{2\pi}V_{Te}} \exp\left(-\frac{v^2}{2V_{Te}^2}\right) \quad (8)$$

The full number of suprathermal electrons is not large about 2% of the thermal electrons.

Let us determine now the transport of energy by fast electrons. Integrating distribution function (7) we can calculate the full power transported out of the turbulent layer by accelerated suprathermal electrons

$$W_a = 2 \int_{4V_{Te}}^{25V_{Te}} \frac{1}{2} m v^2 v f_s dv \approx 2.68 n_0 T_e V_{Te} \approx 1700 \frac{\mu kW}{m^2} \quad (9)$$

We see that the power transported by suprathermal electrons is 20 times larger, than the full power of the pump wave ($W_p \approx 96 \mu kW/m^2$). This result is obviously controversial. It means, that in spite of the fact that mean free path of the suprathermal electrons in ionosphere F-layer is larger than the width of the turbulent region $\lambda_m \approx 180$ m, the electron collisions could not be neglected. The flux of the energy out of the layer is determined by collisions – suprathermal electrons not only leave the layer, but due to collisions they can return back. Thus only the full balance between accelerated inside the layer and returning backward suprathermal electrons can determine correct relation between the pump wave energy losses and acceleration. Just this balance is studied in the theory of multiple acceleration. This theory in modernized form would be considered below.

Note, that the controversial result of electron acceleration in non collisional kinetics shows, that the theory of plasma turbulence saturation should be revised as well: electron collisions outside the Langmuir turbulent layer should be taken into account to organize a correct flow of plasma energy. In other words boundary conditions, which connect turbulent and nonturbulent regions, could affect the process of Langmuir waves saturation inside the turbulent layer.

1.4 Kinetic theory of multiple acceleration

To find the distribution function we have to solve kinetic equation inside and outside acceleration layer and match solutions.

Inside the layer kinetic equation has a form:

$$\frac{\partial f}{\partial t} + v_{\parallel} \frac{\partial f}{\partial z} - \frac{e}{m} E \frac{\partial f}{\partial v_{\parallel}} = 0 \quad (10)$$

Here E is electric field of Langmuir cavitons excited by powerful radio wave

$$E(z, t) = \frac{1}{2} E(z) (e^{-i\omega t} + c.c) \quad (11)$$

The solution of the equation (10) has the form:

$$f(z, \varepsilon_{\parallel}, t) = F(z - Z(t), \varepsilon_{\parallel} - \varepsilon(t)) \quad (12)$$

Here $\varepsilon(t)$, $Z(t)$ the trajectory of a fast particle in electric field E :

$$\frac{dZ}{dt} = V, \quad \frac{dV}{dt} = \frac{e}{m} E, \quad V = \sqrt{2\varepsilon/m}$$

To solve our problem we need to determine the changing of distribution function Δf when a fast particle crossing the acceleration layer. Averaging relation (12) on fast oscillations we find the changing of distribution function after passing through the acceleration layer).

$$\Delta f = \frac{1}{2} \frac{\partial}{\partial \varepsilon_{\parallel}} \left[\langle (\Delta \varepsilon)^2 \rangle \frac{\partial F}{\partial \varepsilon_{\parallel}} \right] \quad (13)$$

According to (11) after averaging on fast oscillations in a fix spatial point energy gain is given by a formulae:

$$\Delta \varepsilon(z, t) = -\Delta \varepsilon \cos \left(\omega t - \frac{\omega z}{v_{\parallel}} + \psi \right) \quad (14)$$

$$\Delta \varepsilon(\varepsilon_{\parallel}) = e |E_k|, \quad E_k = \int_{-\infty}^{+\infty} E(z) e^{-ikz} dz, \quad k = \frac{\omega}{v_{\parallel}} = \omega \sqrt{\frac{m}{2\varepsilon_{\parallel}}}$$

For example, for Gauss like caviton $E(z) = E_0 \exp(-(z/a)^2)$ we have from (11):

$$\Delta \varepsilon = \pi^{1/2} e E_0 a \exp(-ka^2) \quad (15)$$

We see from (15) that only electrons with high enough energy

$$\varepsilon_{\parallel} > \varepsilon_1 = m(\omega a)^2 = T_e \left(\frac{a}{D_e} \right)^2 \quad \left(\frac{a}{D_e} \right)^2 \gg 1 \quad (16)$$

where D_e is a Debay length, could be accelerated. It is easy to see that relation (16) has a universal character and does not depends on a real form of a caviton. It is obvious that if we have some number of noncorrelated cavitons the average energy changing is equal to a sum

$$\langle \Delta \varepsilon^2 \rangle = \sum_i \langle \Delta \varepsilon_i^2 \rangle$$

Now let us consider the solution outside acceleration layer at $z > 0$ and $z < 0$. Kinetic equation has a form:

$$\frac{\partial f}{\partial t} + v \mu \frac{\partial f}{\partial z} = -S \quad (17)$$

Here $f(t, v, \mu, z)$ is the distribution function of electrons, S - collision integral, μ - cosine of an angle between the velocity of electron and z direction. The collision integral consists of two parts: inelastic S_0 and elastic S_1 . The collision integral depends strongly on energy. For energies $\varepsilon < \varepsilon^*$ ($\varepsilon^* \approx 4$ eV) the main collision crosssection is defined by oscillational levels and inelastic collision integral has a diffusive form. But for energies $\varepsilon > \varepsilon^*$ the collision crosssection is determined by optical levels. Thus the collision integral could be presented in a form:

$$S_0(F_0) = -\frac{\partial}{\partial \varepsilon} \left(R \frac{\partial f}{\partial \varepsilon} \right) \quad \varepsilon < \varepsilon^* \\ S_0(F_0) = v_0(\varepsilon) F_0 \quad \varepsilon > \varepsilon^* \quad (18)$$

$$v_0 = v \sigma_0 N_m$$

$$S_1(f_1) = v_1(\varepsilon) f_1, \quad v_0 = v \tilde{\sigma}_i N_m$$

$$\tilde{\sigma}_i = \sigma_i \theta(\varepsilon - \varepsilon^*) + \sum_i \theta(\varepsilon - \varepsilon^*)$$

The σ_0 is the sum of effective crosssection of inelastic collisions with neutral atoms N_m - the total density of neutrals. σ_i is transport electron crosssection of electron elastic collisions

$$\sigma_0 = \sum_k \sigma_{k0} \frac{N_{mk}}{N_m} \quad \sigma_i = \sum_k \sigma_{ik} \frac{N_{mk}}{N_m} \quad N_m = \sum_k N_{mk} \quad (19)$$

The averaged part of total energy of an electron loosing in one collision is small; it is defined by parameter

$$\delta = v_0/v_1 = \sigma_0/\sigma_i \quad (20)$$

High electron energies ($\varepsilon > \varepsilon^*$)

Because of (20) we can expand the equation (17) on small parameter δ . At first approximation the distribution function does not depend on μ . Taking into account first approximation one can find:

$$\frac{\partial F_0}{\partial t} + \frac{v}{3} \frac{\partial f_1}{\partial z} = -S_0, \quad S_0 = \frac{1}{2} \int_{-1}^1 S d\mu \quad (21)$$

$$\frac{\partial F_0}{\partial z} = S_1 = -v_1 f_1$$

Boring in mined the boundary conditions (13) at $z = 0$ in stationary conditions from (21) one can find:

$$\frac{\partial F_0}{\partial z} = -\text{sign}(z) \frac{3v_1}{9v\varepsilon} \frac{\partial}{\partial \varepsilon} \left\{ \int_0^\varepsilon \langle \Delta \varepsilon^2(\varepsilon_{\parallel}) \rangle d\varepsilon_{\parallel} \frac{\partial f_0}{\partial \varepsilon} \right\} \quad (22)$$

It is easy to see that for collision integrals in the form (18) the equation (22) depends on energy parametrically only. So we can rewrite equation (18) in the region of energies

$$\frac{\partial^2 F_0}{\partial \xi^2} = 3\delta F_0, \quad \xi = \int_0^\varepsilon \frac{v}{v_1} dz = \int_0^\varepsilon \sigma_1 N_m(z) dz \quad (23)$$

Taking into account that δ is a slow changing function it is possible to solve (23) in WKB approximation

$$F_0(\varepsilon, z) = F_0(\varepsilon) \left[\frac{\delta(0)}{\delta(z)} \right]^{1/4} \exp \left\{ - \left| \int_0^z \frac{dz}{L_\varepsilon(z)} \right| \right\} \quad (24)$$

$$f_1(\varepsilon, z) = \text{sign}(z) (3\delta)^{1/2} F_0(\varepsilon, z)$$

Here L_ε is a relaxation scale of fast electrons with given energy

$$L_\varepsilon = \frac{v}{(3v_1 v_0)^{1/2}} = \frac{1}{N_m (3\sigma_1 \sigma_0)^{1/2}}$$

Middle electron energies ($T_e < \varepsilon < \varepsilon^*$)

In the region of energies $\varepsilon < \varepsilon^*$ the kinetic equation (21) takes the form:

$$\frac{\partial^2 F_0}{\partial \xi^2} = \frac{\partial}{\partial \varepsilon} \left(R \frac{\partial F_0}{\partial \varepsilon} \right) \quad (25)$$

Since the solution of kinetic equation in the region $\varepsilon > \varepsilon^*$ has a form (24), to match this solution with solution in region $\varepsilon < \varepsilon^*$ we had to search the solution for $\varepsilon < \varepsilon^*$ in a form

$$F_0 = g(\varepsilon) \exp \left\{ - \left| \int_0^z \frac{dz}{L_\varepsilon(z)} \right| \right\}$$

The equation (25) takes the form:

$$\frac{\partial}{\partial \varepsilon} \left(R \frac{\partial g}{\partial \varepsilon} \right) = - \frac{v_0 v_1}{\tilde{v}_1} g$$

Taking into account that

$$v_0 v_1 \ll R \tilde{v}_1$$

One can see from (26) that at first approximation

$$F_0(\varepsilon) = F_0(\varepsilon^*) = \text{const} \quad \varepsilon < \varepsilon^* \quad (27)$$

The distribution function $F_0(\varepsilon) = F_0(\varepsilon, z)|_{z=0}$ according to boundary condition (13) is determined by diffusion equation:

$$\frac{1}{\delta^{1/2} \varepsilon} \frac{d}{d\varepsilon} \left\{ \delta^{1/2} \varepsilon T_{eff}^2 \frac{df_0}{d\varepsilon} \right\} = f_0 - f^{(0)}(\varepsilon), \quad F_0 = f_0 - f^{(0)}(\varepsilon) \quad (28)$$

$$\text{and } F_0(\varepsilon) = F_0(\varepsilon^*) = \text{const}, \quad \varepsilon < \varepsilon^*$$

We introduce here effective temperature :

$$T_{eff}(\varepsilon) = \kappa \bar{\Delta} \varepsilon, \quad \bar{\Delta} \varepsilon = \left\{ \frac{1}{\varepsilon} \int_0^\varepsilon \langle \Delta \varepsilon^2(\varepsilon_{||}) \rangle d\varepsilon_{||} \right\}^{1/2}, \quad \kappa = \frac{3^{1/4}}{2\sqrt{2}} \delta^{-1/4} \quad (29)$$

here $\bar{\Delta} \varepsilon$ is mean electron energy increasing in accelerating layer. The effective temperature increase gradually with increasing energy and saturate at the level.

$$T_{effm} = \kappa(n/2)^{1/2} \Delta \varepsilon_m, \quad \varepsilon_m \leq \varepsilon_1$$

Here n is the number of cavitons in accelerating layer.

The intensity of electron acceleration is defined by parameter

$$\gamma = \frac{T_{eff}}{T}, \quad T = -f^{(0)}(\varepsilon) \left/ \frac{df^{(0)}(\varepsilon)}{d\varepsilon} \right.$$

For weak acceleration ($\gamma < 1$) the distribution function of electrons does not change substantially

$$F_0(\varepsilon) = \left(\frac{T_{effm}}{T} \right)^2 f^{(0)}(\varepsilon)$$

In opposite case of strong acceleration $\gamma \gg 1$ the equation (29) in WKB approximation has the following solution:

$$f_0(\varepsilon) = f^{(0)}(\varepsilon) \quad \varepsilon < \varepsilon_{th}$$

$$f_0(\varepsilon) = f_0(\varepsilon_{th}) \left\{ \frac{\pi \varepsilon_{th} T_{effm} \delta^{1/2}(\varepsilon_{th})}{2 \varepsilon T_{effm}(\varepsilon) \delta^{1/2}(\varepsilon)} \right\}^{1/2} \left\{ \frac{dT_{effm}}{d\varepsilon} \right\}^{-1/2} \exp \left[- \int_{\varepsilon_{th}}^\varepsilon \frac{d\varepsilon}{T_{effm}(\varepsilon)} \right] \quad \varepsilon > \varepsilon_{th} \quad (30)$$

The value ε_{th} in (30) is defined by matching of distribution functions in saddle point

$$T_{effm} = T(\varepsilon_{th}).$$

In the case of strong acceleration in the limit $T_{eff} \approx \text{const}$, $\delta^{1/2} \approx \text{const}$ the distribution function (30) has the form:

$$f(\varepsilon) = CK_0 \left(\frac{\varepsilon}{T_{eff}} \right) \quad (31)$$

Here $K_0(x)$ is a modified Bessel function. The number of fast electrons is defined by the integral

$$N_e^f = \int f(\varepsilon) d^3v$$

But the normalized constant C is determined by the microwave power absorbed by fast electrons. The absorbed power consists of two terms

$$\begin{aligned}
P &= P_1 + P_2 & P_1 &= \frac{16\pi}{3m^2} \left(\delta/3\right)^{1/2} C \varepsilon_{th}^3 K_0 \left(\frac{\varepsilon_{th}}{T_{eff}}\right) \\
P_2 &= \frac{16\pi}{m^2} \left(\delta/3\right)^{1/3} C \int_{\varepsilon_{th}}^{\infty} \varepsilon^2 K_0 \left(\frac{\varepsilon_{th}}{T_{eff}}\right) d\varepsilon
\end{aligned} \tag{32}$$

The first term P_1 in (32) is responsible for losses on rotational levels of molecules and the second one P_2 describes the electron acceleration.

The work will be prolonged as it is formulated in the Introduction. Its final text will be presented in the Final Report.

3. Drift effects on the scale of field-aligned irregularities

Irregularities in the F-region of the ionosphere are strongly elongated along the magnetic field due to dramatic difference in transport coefficients along the magnetic field and transverse to it. Their scales along the magnetic field L_{\parallel} in stationary state were determined previously neglecting plasma drift effects both in linear [Gurevich, 1978 and many other publications] and nonlinear [Gurevich et al., 1995] approximations. One can suppose that drifts could substantially diminish L_{\parallel} for irregularities of sufficiently small transverse (to the magnetic field) scales in the drift direction.

Transport equations along the magnetic field can be written in the form [Gurevich, 1978]

$$-V \frac{\partial T}{\partial x} + \frac{\partial}{\partial z} \left(\kappa_{\parallel} \frac{\partial T}{\partial z} \right) - \delta \nu T + Q = 0 \quad (1)$$

$$-V \frac{\partial N}{\partial x} + \frac{\partial}{\partial z} \left(D_{\parallel} \frac{\partial N}{\partial z} \right) + \frac{\partial}{\partial z} \left(D_{\parallel}^T \frac{\partial T}{\partial z} \right) - \gamma N = 0 \quad (2)$$

Here T and N are disturbances of temperature and density in irregularities, D , D^T , and κ with index \parallel are diffusion, thermal diffusion, and heat conductivity along the magnetic field respectively, γ is the recombination coefficient, δ is the average part of electron energy loss under collisions with ions and neutrals, and ν is the electron collision frequency. In (1),(2) we supposed that irregularities are induced by heating source Q , we neglected all transport across the magnetic field (supposing transverse scale to be not too small) and took into account drift across the magnetic field with the velocity V .

Without drifts longitudinal scale of irregularities is determined by the longitudinal diffusion length $L_N = (D_{\parallel}/\gamma)^{1/2}$ and the longitudinal heat conductivity length $L_T = (\kappa_{\parallel}/\delta \nu)^{1/2}$ [Gurevich, 1978]. In the F-region of the ionosphere $L_T \gg L_N$ [Gurevich, 1978], and longitudinal scale of irregularities is determined mainly by the heat conductivity length L_T . In case of the heating source Q produced by the powerful HF radio waves its scale along the magnetic field is much smaller L_T , and thus its heating can be treated using Dirac delta function $\delta(z)$.

Let us consider heat transfer equation (1). Introducing undimensional variables

$$\tau = \frac{T}{T_0}, \quad \xi = \frac{x}{L_V}, \quad \varsigma = \frac{z}{L_T}, \quad v = \frac{V}{V_0}, \quad q = \frac{Q}{T_0 \sqrt{\kappa \nu}}$$

where $L_V = V_0/\delta \nu$, T_0 is unperturbed electron temperature, and V_0 is a characteristic drift velocity, we can rewrite equation (1) in the following form

$$-v \frac{\partial \tau}{\partial \xi} + \frac{\partial^2 \tau}{\partial \varsigma^2} - \tau + q \delta(\varsigma) = 0 \quad (3)$$

In case of constant drift velocity $v=1$ one can easily found solution of (3) for harmonic in ξ irregularities $\tau \propto \exp(ik\xi)$ as

$$\tau_k = \frac{q_k}{2\sqrt{1+ik}} \exp(-\sqrt{1+ik}\varsigma) \quad (4)$$

where q_k is the respective harmonic component of the heat source. From this solution one can find that irregularities with transversal scale $L < L_V$ in the direction of plasma drift have longitudinal scales decreasing as $(L/L_V)^{1/2}$ as well as the efficiency of their excitation.

In the F-region characteristics scale L_V is of order of kilometer for the drift velocity of order of few ten meters per second.

Above conclusion is valid only if heating source has a fixed position in the ionosphere (say, forming by the radio wave beam). From the other hand small scale (few meters) irregularities – striations [Gurevich et al., 1995] are excited inside the beam, and their source can be considered as drifting with plasma. In this case gradients of the drift velocity becomes essential. Supposing linear gradient of drift velocity along the magnetic field we can rewrite equation (1) using undimensional variables as

$$\frac{\partial \tau}{\partial \xi} \frac{\partial \tau}{\partial \xi} + \frac{\partial^2 \tau}{\partial \xi^2} - \tau + q \delta(\xi) = 0 \quad (5)$$

where characteristic scale L_V is determined now as

$$L_V = \frac{L_T}{\delta v} \frac{dV_0}{dz}$$

Vanishing at infinity solution of (5) for harmonic component $\tau \propto \exp(ik\xi)$ is

$$\tau_k = \frac{(ik)^{-1/3} q_k}{2Ai' \left((ik)^{1/3} \left(\zeta - \frac{i}{k} \right) \right)} \quad (6)$$

Analizing behavior of solution (6) for different values of k one can find that longitudinal scale remains L_T for transversal scales of irregularities $L > L_V$ decreasing for smaller scales approximately as $(2 \lg(L_V/L))^{-1}$. Efficiency of excitation also decreases for small scales as a factor of $k^{-1/3}$.

In the F-region characteristics scale L_V is of order of 10 meters for drift velocity gradient dV_0/dz of the order of several cm/s per kilometer. Thus, we see that drift gradients could essentially affect stationary state of striations of meter transverse scales

In our future work we intend to take into account effects of the drift on the nonlinear structure of striations both already existing in ionosphere and induced by the action of powerful radio wave. The final results would be presented in our Final Report.

References

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